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BEHAVIOR OF A TWO-LEVEL SOLID-STATE MASER

I. R. Senitzky

United States Army Signal Research
and Development Laboratory
Fort Monmouth, New Jersey
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The generation of oscillations in a microwave cavity by means of two-level electron-spin systems has been reported recently by two groups.[1,2] A characteristic of both results is the amplitude modulation of the power from the cavity during the time of oscillation. This modulation is quite regular, has a well defined period, and, in one of the results,[2] is greater than fifty percent. The explanation of this modulation is not given with certainty, but the authors suggest that perhaps field inhomogeneities are responsible through the interference between different spin packets. However, some of the authors have since concluded that this explanation cannot be a valid one for modulation of such regularity and reproducibility,[3,4] and there seems to have been aroused considerable interest in the cause of the modulation. The purpose of this letter is to suggest an explanation. It will be shown, under simplifying assumptions, that when a system of spins is in resonance with the cavity, and when damping effects caused by relaxation phenomena and cavity losses are not excessive, modulation of the cavity power output is exactly what is to be expected on the basis of a dynamical analysis.

Consider a cavity inside of which there is a crystal containing the spin systems. In order to exhibit clearly and simply the reason for the modulation, we neglect complicating features which are not essentially related to the problem. We will thus ignore, at first, relaxation effects and cavity losses, which will be discussed briefly later, qualitatively. Also, a classical treatment of the total angular momentum of the spin systems and of the electromagnetic field inside the cavity will suffice.

We take the dc magnetic field, H_0, to lie along the z axis. The rf magnetic field is denoted by

$$\vec{H} = Q(t)\,\nabla\times\vec{u}(\vec{r}),$$

where the function $\vec{u}(\vec{r})$ describes the spatial variation of the field for the mode with which we are concerned, and is normalized over the volume of the cavity; that is, $\int_V u^2\,d^3\vec{r} = 1$.

We consider the case in which the rf field at the crystal is perpendicular to the dc field; and, for simplicity, we assume that the rf field is uniform over the volume of the crystal. We can, therefore, consider it to lie along the x axis, so that, at the crystal, $\nabla\times\vec{u}(\vec{r}) = \hat{x}\,a_x$, where we have set $\vec{a} = \nabla\times\vec{u}$. The total angular momentum of the spin systems will be denoted by $\vec{S}$, and the corresponding magnetic moment by $-\alpha\vec{S}$.

The equations of motion for Q and S_x are

$$\ddot{Q} + \omega^2 Q = -\,4\pi C^2 \alpha\, a_x\, S_x, \tag{1}$$

$$\ddot{S}_x + \omega_0 S_x = \omega_0\,\alpha\, a_x\, S_z\, Q, \tag{2}$$

where ω is the resonant frequency of the cavity and $\omega_0 = \alpha H_0$. For the case of resonance between the spin system and the cavity, which is the case we are considering, $\omega = \omega_0$. Now, S_z varies much

more slowly than S_x or Q. This may be shown analytically from a more complete set of equations of motion (which would include an equation for S_z), or may be seen from a picture of the precessing spin system. Moreover, while Q and S_x oscillate through zero with approximately sinusoidal motion, S_z need not change by much compared to its average value. The value of S_z is determined by the total energy of the spin systems, which is $\alpha S_z H_0$.

Consider now a situation in which the total energy in the cavity (that of the spins and the rf field) is less than half the maximum spin energy. Then, S_z is negative. Replace S_z in Eq. (2) by its average value, $\bar{S}_z$. Eqs. (1) and (2) now become the linear equations for two coupled harmonic oscillators, the field and the spins. It is well known, and can also be calculated easily from the equations of motion, that in such a situation the energy transfers back and forth from one oscillator to the other, the modulation frequency of the energy of each oscillator being ν_e, where $\nu_e^2 = -\pi^{-1}\omega^{-1}C^2\alpha^2 a_x^2 \bar{S}_z$. The accuracy of this expression increases as the total energy is reduced, since then $\bar{S}_z$ is a better approximation for S_z, and the coupling between the two oscillators becomes more linear. But it seems evident that the explanation for the modulation is the well known periodic transfer of energy between two coupled oscillators.[5] Damping caused by relaxation effects and cavity losses may destroy the modulation only if it is large enough to dissipate all the energy in one modulation cycle. Otherwise, it will affect the modulation phenomenon mainly through the (slight) dependence of the modulation frequency on the total energy, because of the nonlinear coupling. Nonuniformities in the dc field might conceivably blur the modulation phenomenon; and, according to the above theory, it would seem that the more uniform the magnetic field is, the more clear-cut the modulation would be.

Finally, we compare the calculated modulation frequency ν_e, above, with the experimental frequency, which can be obtained easily from the display of power output vs time in Fig. 1 of reference 2. For the lower-energy portion of the curve it is about 2×10^6 cycles per second. The authors also give the following data: $\omega/2\pi = 9\times10^9$ cycles/sec, the total number of electron spins is $\sim10^{18}$, and the temperature of the crystal is 4.2°K. The last two facts imply that $\bar{S}_z \sim 10^{16}\hbar$. For a_x^2 we assume the value of $|\vec{a}|^2$ averaged over the volume V of the cavity, $\langle a^2\rangle_{Av} = \omega^2 C^{-2} V^{-1}$, where V is 12 cm³.[3] Bearing in mind that $\alpha = 2\mu_0/\hbar$ where μ_0 is a Bohr magneton, we obtain $\nu_e = 2\times10^6$ cycles/sec, in agreement with the experimental data.

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[3]Chester, Wagner, and Castle, Solid-State Maser Symposium, U. S. Army Signal Research and Development Laboratory, Fort Monmouth, New Jersey, June 12, 1958 (to be published).

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DIRECT OBSERVATION OF SPIN-WAVE RESONANCE*

M. H. Seavey, Jr., and P. E. Tannenwald†

Lincoln Laboratory,
Massachusetts Institute of Technology,
Lexington, Massachusetts
(Received August 11, 1958)

It has been generally understood that the effects of exchange interactions on ferromagnetic resonance in metals can be observed experimentally provided the rf field gradient is sufficiently great.[1-6] A theoretical treatment for the case of a thin metallic film leads to the prediction of "body resonances" due solely to such an rf gradient.[7,8] However, Kittel[9] has now proposed a boundary condition which permits excitation of these resonances even by a uniform rf field and he assigns to them the more appropriate term "spin-wave resonances." Here the spins are simply assumed to be pinned down at the boundaries of the specimen due to a surface anisotro-

py. Kittel points out that thin films should be particularly well suited for observation of spin-wave resonance.

Figure 1 shows the microwave absorption spectrum at 8890 Mc/sec for a 80% - 20% Permalloy film of 5600 A thickness. The dc field is perpendicular to the film surface: The various

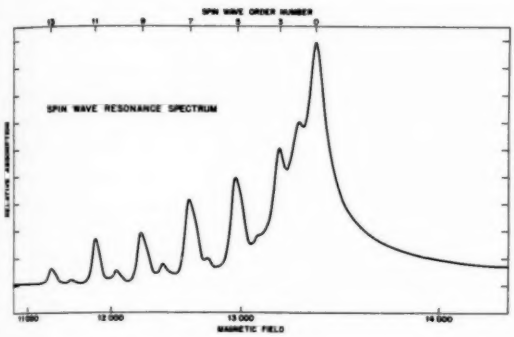

FIG. 1. Experimental observation of spin wave resonances in a 5600 A film of permalloy.

resonances are interpreted as arising from excitation of spin waves obeying the dispersion law

$$\omega = (2A\gamma/M)\,k^2 + \omega_0, \qquad (1)$$

and the Kittel condition for spin wave resonance

$$k = p\pi/L, \qquad (2)$$

where A is the exchange coupling constant, M the magnetization, L the film thickness and p an integer. A value for A of $(2.0 \pm 0.03) \times 10^{-6}$ erg/cm is obtained by fitting the resonance field data to the above expressions (1) and (2) for the odd modes $p = 7$ to 13. The small peaks between the large ones in the figure may be interpreted as even-mode resonance (even number of half-wavelengths) which in the absence of an rf field gradient would not be excited.

The exact exchange-effect theory worked out by the authors[7] when applied to this particular film yields highly damped peaks at the field values predicted by (1) and (2) for $p > 3$. The large magnitudes which are actually observed can thus be explained on the basis of the Kittel boundary condition, which permits odd-mode excitation even by a uniform field. The exact calculation,[7] however, shows that for $p < 3$ the dispersion relation (1) does not hold and the Q (proportional to $\mathrm{Re}k/\mathrm{Im}k$) of the spin waves[10] becomes small. In fact the mode $p = 1$ can be shown to fall on the high-field side of the main resonance and to be damped out in a distance less than one-half wavelength. Hence the first spin-wave resonance in the figure corresponds to $p = 2$. It should also be pointed out that the exact theory yields a line width for the main resonance of ~150 oe, a value quite close to the observed one.

The resonance field data only approximately follow the p^2 law for $p > 3$; a better experimental fit could be obtained by using an exponent less than two. This may be explained qualitatively by considering the strength of the spin pinning. It can be seen from Eq. (15) of reference 9 that the pinning at the boundaries gets weaker with increasing spin wave number thus causing a progressive shift of the resonance peaks.

The direct observation of spin-wave resonances thus provides a straightforward way of obtaining the exchange constant A from the location of the spin-wave peaks. Furthermore, the damping of isolated spin waves of a given wave number is obtained simply by measuring the width of the peaks. The nature of the surface anisotropy may be evidenced by deviations from the p^2 law.

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† Staff Members.

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ELECTRON-NUCLEAR DOUBLE RESONANCE OF F CENTERS IN LITHIUM FLUORIDE*

Norman W. Lord
Applied Physics Laboratory,
The Johns Hopkins University,
Silver Spring, Maryland
(Received July 7, 1958)

In LiF, the spin resonance of the F-center electron is partially resolved by its hyperfine interactions with neighboring nuclei.[1] Discrepancies between observed and theoretically predicted values of these interactions[2] are most serious for the lighter alkali halides. This situation motivated a more detailed experimental examination of the electronic structure of this center using the new double resonance technique.[3]

A crystal of LiF was heavily irradiated with 50-kv x-rays thereby inducing a high concentration of F centers. The sample, on a long rod, was placed in thermal contact with a metal container filled with liquid helium, at the same time projecting into a rectangular X-band cavity. While partially saturating the electron spin resonance with a microwave field at 9200 Mc/sec, an additional intense radio-frequency field was superimposed on the sample using fine wires suspended inside the metal cavity. As the frequency of this auxiliary field is varied, the spin resonance signal changes whenever the field stimulates nuclear spin transitions between the Zeeman hyperfine levels of the F-center electron. The change becomes appreciable as this hyperfine interaction "resonance" of the neighboring nucleus is saturated. A maximum steady state change is approached monotonically by saturating both the electron and nuclear resonances. The rf field intensity for this was around 0.5 gauss and depended somewhat on the particular nucleus involved.

Both the microwave and radio-frequency magnetic fields were perpendicular to a steady magnetic field, H, of approximately 3200 oersteds. The crystal [100] axis was along the length of the bar and perpendicular to H. On rotation of the sample about this axis, H remained in a (100) plane and varied its angle with respect to a [100] axis. The allowed nuclear spin transitions are $\Delta m_\alpha = \pm 1$. The separations between the connected spin levels of nucleus α are

$$h\nu_\alpha = |\tfrac{1}{2}\ [E_\alpha{}^S + E_\alpha{}^{\text{dipole}}]\ \pm \mu_\alpha H/I_\alpha| \ , \quad (1)$$

in the notation of reference 2. For analysis of the experimental data, the dipolar part can be re-expressed as

$$E_\alpha{}^{\text{dipole}} = (3\cos^2\theta_\alpha - 1)\ 2\beta\ \frac{\mu_\alpha}{I_\alpha} \int \varphi_n(n) \times \frac{[3\cos^2\gamma_{\alpha n} - 1]}{r_{\alpha n}{}^3}\ \varphi_n{}^*(n)\ d\tau_n \ , \quad (2)$$

where θ_α is the angle between H and the line, α-n, joining the F center, n, to nucleus α, while $\gamma_{\alpha n}$ is the polar angle of integration measured from this axis.

These interactions of (1) give rise to an angle-dependent double resonance spectrum which can be grouped into sets of lines associated with the angle-independent term $E_\alpha{}^S$. The nuclei responsible for any set are crystallographically equivalent with respect to the halide ion vacancy of the F center, and form a shell of sites whose distance from this vacancy is $\sqrt{N}$ times the lattice parameter, N being integral. Lines were observed for shells out to $N = 6$. Figure 1 shows the sets of double-resonance lines for the first 3 shells surrounding the F center when the direction of H in a (100) plane is varied between 0 and $\pi/4$ with respect to a [100] axis.

The experimental results for the constants, $E_\alpha{}^S/h$ and $E_\alpha{}^{\text{dipole}}/h(3\cos^2\theta_\alpha - 1)$, which are direct measures of the F-center electronic wave function φ, are presented in Table I. The line widths of the double-resonance lines are also included as a measure of the errors inherent in these measurements.

The contact term result for shell 1 agrees very well with its earlier determination. The decline of this term in shell 2 is much faster than expected either from the theory[2] or experiment,[1] suggesting an even greater concentration of φ within the vacancy. However, the earlier results rested on the apparently unjustified as-

Table I. Experimental values of hyperfine interaction energies in megacycles per second for the first three shells of nuclei surrounding the F center in LiF. The last column contains the half-width of the double-resonance spectral lines associated with the shell.

Shell	$\frac{E_\alpha{}^S}{h}$	$\frac{E_\alpha{}^{\text{dipole}}}{h(3\cos^2\theta - 1)}$	$\Delta\nu_\alpha$
1	38.5	+3.1	0.48
2	1.34	+0.56	0.07
3	0.54	+0.64	0.03

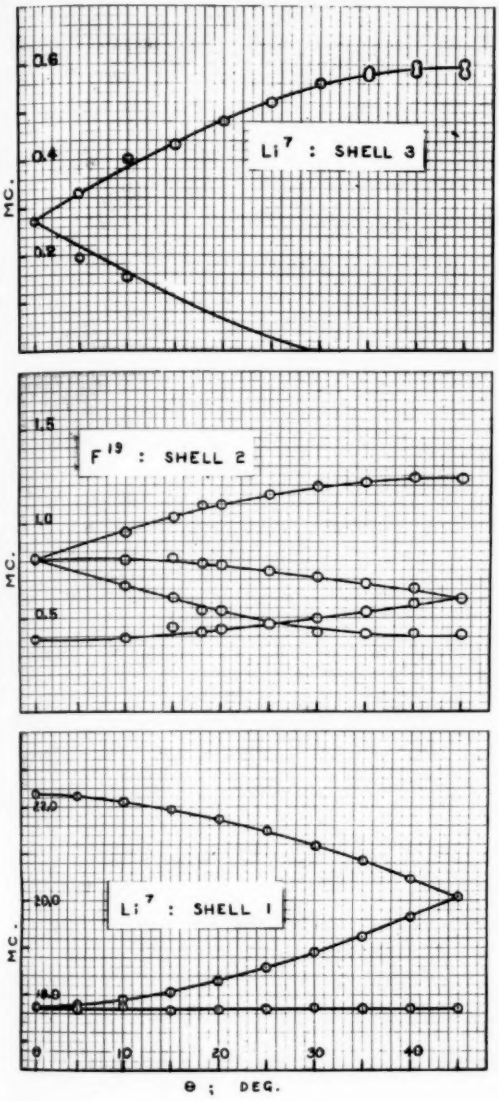

FIG. 1. Variation of double-resonance spectra as a function of the angle θ between H and the [100] axis in the (010) plane. Splitting in shell 3 is due to slight elevation of H out of this plane.

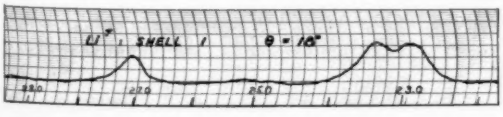

FIG. 2. Portion of double-resonance spectrum due to the first shell, demonstrating the extraordinary line width.

sumption that the width of each of the 19 partially resolved components was predominantly the hyperfine interaction of the F^{19} nuclei of shell 2. As a contrast, the dipolar interaction shows a remarkable persistence. This fact alone made portions of the outer shell spectra observable, their E_α^S being very much smaller.

It was expected that the lack of axial symmetry of φ about α-n axis would be much stronger in the more closely packed LiF cyrstal than in KCl. No evidence of this appeared either as a quadrupole splitting of shell 1 and shell 3 lines, or as an additional θ-dependence of shell 2 lines. Appreciable axial asymmetry would cause the intersecting low-frequency traces in Fig. 1 to be different at 0° and 45°, instead of being practically equal. Work in progress on the intermediate alkali halides may clarify this point.

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3 G. Feher, Phys. Rev. 105, 1122 (1957).

"RADIATION BELT" AND TRAPPED COSMIC-RAY ALBEDO*

S. F. Singer

Department of Physics, University of Maryland, College Park, Maryland

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The existence of a belt of unidentified ionizing radiation observed by the Explorer satellites[1] has led to speculations concerning the nature and source of this radiation. In a recent Letter Dessler[2] suggests that the radiation is related to auroral particles and is accelerated by hydromagnetic waves. This suggestion runs into several difficulties: (i) Hydromagnetic waves traveling along lines of force are probably the accelerating agent for auroral protons,[3] but only in the auroral zone (60°-70° magnetic latitude), where both waves and particles are trapped and can exchange energy.[4] (ii) Waves probably do not propagate perpendicular to the lines of force because of the strong reduction in conductivity

due to a large neutral atom component in the exosphere.[4] (iii) In any case such waves could not accelerate particles very effectively at the equator and at the same time keep the particles trapped. (iv) The relative constancy of the observed radiation belt is difficult to reconcile with auroral phenomena, or indeed any solar corpuscular activity, and suggests instead a constant source for at least part of the radiation. (The observations certainly do not exclude a 27-day recurrence; furthermore auroral radiation might somehow diffuse to the equator.)

We have therefore been tempted to apply some earlier work[5] on cosmic-ray albedo, and particularly on albedo trapped in the geomagnetic field, to see how far it accounts for the satellite observations. Irrespective of such observations, however, the existence of trapped albedo is a necessary consequence of the impact of primary cosmic rays on the earth. Calculations (previously unpublished) were therefore started in connection with cosmic-ray experiments planned for the Farside Rocket (a 4000-mile Air Force rocket eventually fired in the Pacific in November, 1957).

The trapping of charged particles in the geomagnetic field, their diffusion, and eventual leakage has been studied in great detail for low-energy (~20 kev) solar corpuscular radiation. The azimuthal drift of these trapped protons leads to a Störmer ring current which is responsible for magnetic storms.[4] A fraction of the particles is accelerated to auroral energies (~500 kev) by hydromagnetic waves.[3]

We can apply similar considerations to trapped cosmic-ray albedo. We first assume (and justify it below) that the particles have a radius of curvature

$$\rho = (pc/e)\,(1/300\,B) \ll B/\mathrm{grad}\,B = r/3, \qquad (1)$$

so that their magnetic moment $\mu = \frac{1}{2} m v_\perp^2/B$ is a constant of the motion, and that therefore[6]

$$\sin^2\alpha/B = \sin^2\alpha_0/B_0 = \text{constant}. \qquad (2)$$

B_0 and α_0 refer to the field and the particle's pitch angle in the equatorial plane. If the line of force intersects a given altitude level at latitude λ, then the range of pitch angles of particles which stay trapped above this level is

$$\pi/2 \geq \alpha_0 \geq \alpha_C = \sin^{-1}[\cos^3\lambda\,(1+3\sin^2\lambda)^{-\frac{1}{4}}]. \qquad (3)$$

We can now calculate the mean lifetime in the trapping region, as well as the leakage into the lower atmosphere where the particles are assumed to die. We also calculate the pitch angle distribution $n(\alpha_0)\,d\alpha_0$ giving the concentration of particles per cm^3 in $d\alpha_0$ at α_0. Let

$$(\Delta\alpha)^2 = Dt, \qquad (4)$$

where D is a diffusion coefficient in pitch angle. Assume an isotropic source function

$$q(\alpha_0) = Q\sin\alpha_0\,. \qquad (5)$$

Then we can set up a (steady-state) diffusion equation in α:

$$\partial n/\partial t = \mathrm{div}\,j + q = 0. \qquad (6)$$

Here

$$j = D\,\partial n/\partial\alpha, \text{ and } \mathrm{div}\,j = \frac{\partial}{\partial\alpha}\left(D\frac{\partial n}{\partial\alpha}\right) = D\frac{\partial^2 n}{\partial\alpha^2}. \qquad (7)$$

We finally have

$$D\,\partial^2 n/\partial\alpha^2 + Q\sin\alpha = 0. \qquad (8)$$

Introducing the boundary conditions

$$\partial n/\partial\alpha = 0 \text{ at } \alpha = \frac{\pi}{2};\ n(\alpha_C) = 0, \qquad (9)$$

and integrating, we find the pitch angle distribution at the equator (see Fig. 1):

$$n(\alpha_0) = (\sin\alpha_0 - \sin\alpha_C)\,Q/D. \qquad (10)$$

The rate of leakage of particles is $j(\alpha_C) = Q\cos\alpha_C$ and this equals the total production rate. The total concentration of particles N(per cm^3) is given by

$$N = 2\int_{\alpha_C}^{\pi/2} n(\alpha)\,d\alpha$$

$$= 2\,[\cos\alpha_C - (\sin\alpha_C)\,(\pi/2 - \alpha_C)]\,Q/D, \qquad (11)$$

and the mean lifetime T is determined as

$$T = N/Q = 2\,[\cos\alpha_C - \sin\alpha_C(\pi/2 - \alpha_C)]/D. \qquad (12)$$

The function in square brackets in (12) increases from 0 at $\lambda = 0$ to 1.0 at $\lambda = 90°$. Its variation can generally be neglected in relation to D and Q; it can be taken ~ 0.1.

In the simple derivation above we have neglected the variation of D due to a difference in atmospheric density along the particle's trajectory. Since $\alpha_0 = \alpha_C$ corresponds to particles reaching the lowest point of reflection, they will be most affected by the higher density. Its effect on $n(\alpha_0)$ is indicated in Fig. 1 by the dashed curve. Magnetic bremsstrahlung loss and catastrophic removal of particles are discussed elsewhere[7] and shown to be small.

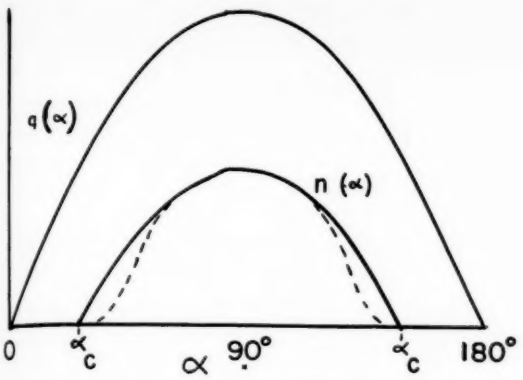

FIG. 1. The assumed isotropic source function $q(\alpha_0)$ and the resultant angular distribution of trapped particles $n(\alpha_0)$ plotted against the particles' pitch angle when they cross the plane of the magnetic equator. Note the absence of particles with $\alpha_0 < \alpha_C$ due to escape from the trapping field. The dashed curve indicates the effect on $n(\alpha_0)$ of the variation of atmospheric density with altitude.

The diffusion coefficient D used above has been evaluated for the mixture of atoms and ions present in the exosphere, using the theory for scattering of relativistic particles under screened and unscreened conditions.[7] Appropriate mean values for the inner exosphere and outer exosphere (where hydrogen predominates) are

$$D_{in} = 20\, cd/\gamma^2; \quad D_{out} = 10\, cd/\gamma^2, \qquad (13)$$

where d is the density in g/cm^3, and $\gamma^2 = (1 - \beta^2)^{-1}$. The transition altitude (~1000 km) should therefore show itself as a sharp increase in N, assuming other quantities in (11) to be constant. If Q is independent of d (or if its dependence is known), then the increase of N with altitude can be used to derive the variation of density with altitude, thereby the scale height, and hence the kinetic temperature of the base of the exosphere.

*Supported by the Air Force Office of Scientific Research.

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CORRELATION OF COSMIC-RAY INTENSITY AND SOLAR ACTIVITY

H. V. Neher

California Institute of Technology, Pasadena, California

and

S. E. Forbush

Carnegie Institution of Washington, Washington, D. C.

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The present International Geophysical Year was chosen to include the most likely period of maximum activity of the sun. It is probably too early to tell whether or not the maximum of the current cycle has yet been reached, but it is already certain that the yearly average of the Zurich sunspot numbers for 1957 is much higher than ever before observed.[1] It is therefore of interest to see what has been the effect on cosmic rays.

In analyzing the data for long periods of time from the Carnegie Institution ionization chambers, Forbush[2] in 1954 found an inverse relationship between solar activity, as measured by Zurich sunspot numbers, and cosmic-ray intensity. Also Neher and Forbush[3] showed in 1952 that there was a good correlation for at least a few weeks between the ionization due to cosmic rays at balloon altitudes at geomagnetic latitude 56°N, the ionization at ground level at Cheltenham and Huancayo, and the neutron intensity at Sacramento Peak, New Mexico, and Climax, Colorado.

It is the purpose of this letter to point out the following relations: (a) The yearly averages of the ionization data at Huancayo correlate very well with the average value of the ionization measured at 90 000 ft, or 15 g cm^{-2} at Thule, Greenland. These latter values were made over about a 2-3 week period during the month of August of the particular year.[4] (b) There is also a very good anti-correlation with solar activity

as measured by the yearly average of Zurich sunspot numbers for the same period. These relationships are shown in Fig. 1.

The large ratio of 19 to 1 for the percentage change near the north geomagnetic pole to that at the equator is due primarily to the large numbers of low-energy particles in the primary radiation which can get through the earth's magnetic field at Thule and penetrate 15 g cm^{-2} of air and which were present in some numbers during the solar minimum of 1954.

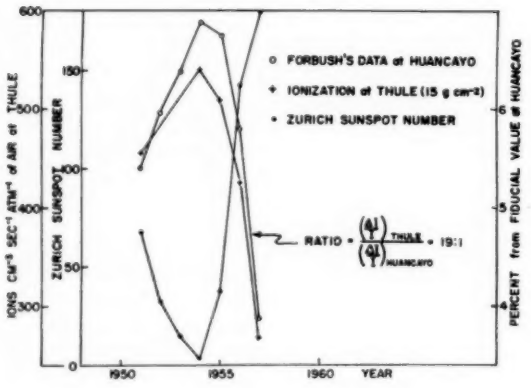

FIG. 1. Long-term correlation between ionization due to cosmic rays at high altitudes (15 g cm^{-2} pressure) near the north geomagnetic pole, the ionization at Huancayo, Peru, and the Zurich sunspot numbers.

The above therefore constitutes further evidence that for these long-time effects, (a) the changes are world wide, (b) the low energy particles are affected more than those of higher energy, (c) the average, yearly Zurich sunspot numbers are a good index of the long-term effect of the sun on the intensity of cosmic rays as measured on the earth.

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PHOTOPION CROSS SECTIONS AND A SECOND RESONANCE*

Ronald F. Peierls
Laboratory of Nuclear Studies,
Cornell University,
Ithaca, New York
(Received August 1, 1958)

In a recent letter[1] Wilson has suggested that the rise in the photopion cross sections above 500 Mev may be interpreted as being due to another resonance in the pion nucleon system. Since the π^+ cross section is the larger, the resonance must presumably be in a $T = 1/2$ state. If we examine the observed angular distributions, it is also possible to determine the probable angular momentum and parity of such a state.

We consider a scheme indicated roughly in Fig. 1. There are three important contributions to the photopion cross section up to about 900 Mev: (A) the $J=3/2$, $T = 3/2$, p-wave resonance at about 300 Mev; (B) the proposed $T=1/2$ resonance at about 700 Mev; (C) the "direct photoelectric" production (s-wave, electric dipole), occurring only for π^+. If A, B, C are the three corresponding complex amplitudes, then for the total cross sections we have

$$\sigma(\pi^+) \propto \frac{2}{3}|A|^2 + \frac{4}{3}|B|^2 + |C|^2,$$
$$\sigma(\pi^0) \propto \frac{4}{3}|A|^2 + \frac{2}{3}|B|^2. \qquad (1)$$

At the peak,~ 700 Mev, one has $\sigma(\pi^+) \cong 2\sigma(\pi^0)$ whence

$$|C|^2 \cong 2|A|^2 \qquad (2)$$

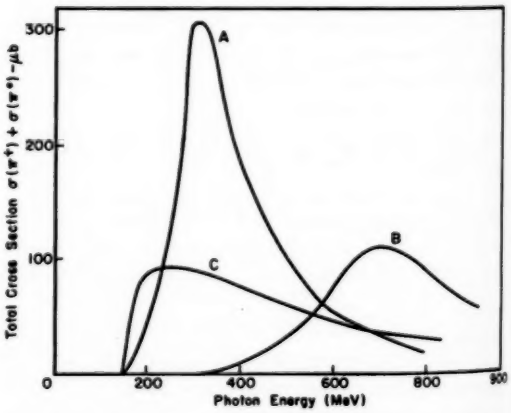

FIG. 1. Major contributions to photoproduction of pions below 900 Mev.

For energies below 350 Mev, the main features of the observed distribution can be explained by A and C alone. For π^0 there is a $2 + 3\sin^2\theta$ distribution throughout, while the π^+ has one half of this together with an isotropic component and an interference term

$$-2\left(\frac{1}{3}\right)^{\frac{1}{2}} \mathrm{Re}(A^*C)\cos\theta\,.$$

The phase of C remains small; that of A increases through 90° at resonance, so that this term gives an asymmetric contribution having a backward maximum below resonance and a forward maximum above. For the measurements near 700 Mev the results are somewhat similar[2-4]: the π^0 shows a steady symmetric distribution with a maximum at 90°, while the π^+ distribution has a strong asymmetry which, however, is peaked forward below, as well as above, the resonance. The π^0 distribution rules out $J=5/2$ for the state B (which would require a dip at 90°) and is quite consistent with $J=3/2$. However, the asymmetric term in the π^+ distribution cannot be interpreted as interference between B and C since it does not change sign on going through resonance. We must conclude that the interference between B and C is symmetric about 90° and thus that B and C have the same parity. Hence the most likely assignment for the proposed resonance is $J=3/2$, odd parity ($D_{3/2}$).

This assignment has one difficulty: If A and B are assigned opposite parity, then there must be an interference term (also $\cos\theta$ in this case) between A and B which should also show up in the π^0 distribution. However, the coefficient of this term involves $-\mathrm{Re}(A^*B)$. Since the widths of the resonances A and B are roughly comparable with the separation between them, the phase of B is small at the maximum of A and increases to 90° at 700 Mev, while the phase of A increases from 90° to near 180°. Hence the relative phase of A and B remains near 90° throughout this region and thus any interference term will be suppressed. Above 700 Mev, this relative phase decreases from 90°, and we expect a backward-peaked distribution to occur. This is, in fact, observed at these energies.

It still remains to explain the observed asymmetry in the π^+. As we have seen, the interference between A and C above the maximum of A is peaked forward, and this asymmetry would not be affected by B going through resonance. At the second resonance the π^+ angular distribution may be written

$$|C|^2 + \frac{1}{6}(2+3\sin^2\theta)(|A|^2+2|B|^2) - 2\left(\frac{1}{3}\right)^{\frac{1}{2}}\mathrm{Re}(C^*A).$$

(At resonance, B^*A and B^*C are purely imaginary.) Since A and C are about 180° out of phase at this energy, and because of (2), this becomes

$$\left(2+2\left[\frac{2}{3}\right]^{\frac{1}{2}}\cos\theta\right)|A|^2 + \frac{1}{6}(2+3\sin^2\theta)(|A|^2+2|B|^2).$$

This agrees with the shape of the actual distribution. Since the observed ratio of maximum to minimum is about 3:1, it follows that

$$|B|^2 \cong 4|A|^2 = 2|C|^2.$$

This determination of the relative magnitudes of A, B, C from the data at one energy, assuming a $D_{3/2}$ resonance, agrees extremely well with the extrapolation of lower energy data for A and C — certainly within the possible accuracy of such an oversimplified calculation.

In pion scattering experiments, a $T=1/2$ peak is also observed, but spread over a much wider range of energies. This can be interpreted as two overlapping levels[1] of which the lower corresponds to that seen in photoproduction. We may then ask why the higher one is not also seen, or at at least does not overlap so strongly. From the magnitude of the total cross sections this level must have $J \geq 5/2$, and thus may also be a D-state resonance. However, in photoproduction a $D_{3/2}$ level can be excited by electric dipole while a $D_{5/2}$ state requires magnetic quadrupole. Hence, as compared with scattering, the $D_{5/2}$ state should be much less important.

We may conclude that, assuming the existence of a resonance, the angular distributions greatly restrict its possible nature. The fact that we are then led uniquely to reasonable values for the ratios $A{:}B{:}C$, consistent with a resonance picture, seems in itself some evidence of the validity of the assumption.

The author is greatly indebted to Professor R. R. Wilson and Professor H. A. Bethe for much stimulating discussion and valuable suggestions.

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*Work supported by the joint program of the Office of Naval Research and the U. S. Atomic Energy Commission.

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ELECTRON POLARIZATION IN THE β-DECAY OF RaE†

Wolfgang Bühring and Joachim Heintze

Zweites Physikalisches Institut
der Universität Heidelberg
Heidelberg, Germany
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The β-decay of RaE(Bi^{210}) is a first-forbidden transition ($1^- \rightarrow 0^+$). The energy spectrum of a nonunique first-forbidden transition usually has the statistical shape and the electron polarization is $P=-v/c$, as in allowed transitions.[1] It is well known,[2] however, that the spectrum of RaE departs from the statistical shape substantially. Due to a particular relation between the nuclear matrix elements, the normally dominating energy-independent Coulomb terms are small and the energy-dependent terms come into play. Consequently one expects in this case the electron polarization to deviate from the normal value $P=-v/c$. Indeed, Geiger et al.[3] have found in a Moeller scattering experiment a smaller effect for RaE than for some other β-emitters investigated.

We have measured the polarization of RaE electrons using a method described in detail in previous publications,[4] which is suitable for accurate relative measurements of the electron polarization of two β emitters having similar energy distributions. The instrument is sensitive for electrons in an energy range from ~250 kev to ~600 kev. In this method the electrons are first deflected by multiple scattering with a thick Cu scatterer and then by single scattering at a thin Pt or Au foil. The scattering angles are 90° and 135°, respectively. The multiple scattering turns the original longitudinal polarization of the electrons into a polarization having a well defined transversal component, which in turn is measured by the azimuthal asymmetry of the single scattering. The asymmetry x is obtained from the ratio of the counting rates L (left) and R (right) of two β counters mounted at 135° with respect to the beam incident on the single-scattering foil: $L/R=\alpha(1-x)/(1+x)$. The factor α, arising from the difference in counter efficiencies, can be eliminated by interchanging the counters, a procedure easily possible in our arrangement. The relationship between x and the polarization can be written in the following form[5]:

$$x=-f_N\frac{\int (Pc/v)N(E)\,w(E)\,dE}{\int N(E)\,w(E)\,dE}=-f_N\left\langle\frac{P}{v/c}\right\rangle_{\mathrm{Av}} \tag{1}$$

$N(E)$ being the β spectrum and $w(E)$ the total efficiency of the apparatus for electrons of energy E. The factor f_N depends on the energy distribution of the β particles. The essential point is that this dependence is small in our experimental arrangement. This is the case even for comparatively thick single-scattering foils. A foil of 2.03-mg/cm² Pt has turned out to be suitable.

We have compared the experimental asymmetry x of RaE ($E_{\max}=1.15$ Mev) with that of Tl^{204} ($E_{\max}=0.76$ Mev) and Y^{91} ($E_{\max}=1.54$ Mev). Tl^{204} and Y^{91} are unique first-forbidden transitions and therefore have the polarization $P=-v/c$.[1] The results are as follows:

$$x(Tl^{204})=f_N(Tl^{204})=(6.42\pm0.05)\%,$$

$$\frac{x(Y^{91})}{x(Tl^{204})}=\frac{f_N(Y^{91})}{f_N(Tl^{204})}=1.016\pm0.014,$$

$$\frac{x(\mathrm{RaE})}{x(Tl^{204})}=\frac{f_N(\mathrm{RaE})}{f_N(Tl^{204})}\left\langle\left(-\frac{P}{v/c}\right)_{\mathrm{RaE}}\right\rangle_{\mathrm{Av}}$$

$$=0.832\pm0.018.$$

These data are mean values of two independent statistically consistent sets of measurements with two RaE sources differing in their effective thickness by a factor of 3. The following corrections have been applied:

(1) Corrections for the instrumental asymmetry. This asymmetry has been measured and the correction shifts the final result by 2%.

(2) Corrections for the depolarization in the source, based on experimental and theoretical results. In the most unfavorable case, this correction amounts to $(3\pm1.5)\%$.

It remains to consider the influence of the β spectra on f_N. Experimentally we have obtained $f_N(Tl^{204})=f_N(Y^{91})$. We have investigated in detail the energy dependence of f_N. It turns out that within an error of <1%, we have $f_N(\mathrm{RaE})/f_N(Tl^{204})=1$. The electron distributions, $N(E)\times w(E)$, of Tl^{204} and RaE are much more alike than those of Tl^{204} and Y^{91}, as can be seen from Fig. 1. We obtain, for the average electron polariza-

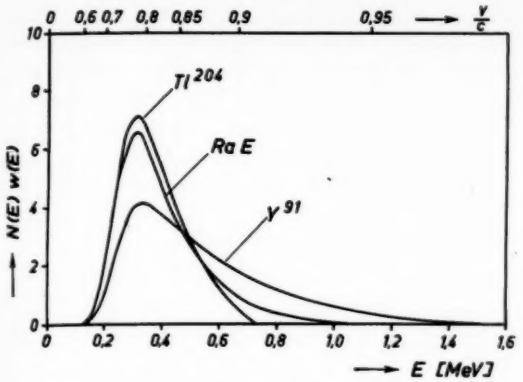

FIG. 1. The electron distributions $N(E)w(E)$ are plotted in arbitrary units. $N(E)$ is the β spectrum, $w(E)$ the total efficiency of the apparatus for electrons of energy E. $w(E)$ has been calculated approximately.[4]

tion defined in Eq. (1) for RaE,[6]

$$\left\langle -\left(\frac{P}{v/c}\right)_{\mathrm{RaE}}\right\rangle_{\mathrm{Av}} = 0.83 \pm 0.02.$$

RaE is the first example of a β emitter which has an electron polarization clearly deviating from $P=-v/c$. Curtis and Lewis[7] have calculated formulas for the polarization and the spectrum shape factor of RaE using STP coupling. For VA coupling[8] the formulas are easily obtained by changing the nuclear matrix elements and coupling constants and reversing the sign of the neutrino momentum,[5] if the two-component theory of the neutrino is valid. The β decay of RaE is governed by three nuclear matrix elements, $\int\vec{\sigma}\times\vec{r}$, $\int\vec{r}$, $\int\vec{\alpha}$. It is convenient to introduce two real[9,10] parameters ξ_1 and η_1 by the relations

$$\xi_1 C_A \int\vec{\sigma}\times\vec{r} = iC_V\int\vec{r},$$
$$\eta_1 C_A \int\vec{\sigma}\times\vec{r} = \left(\frac{\alpha Z}{2\rho}\right)^{-1} C_V \int\vec{\alpha}, \qquad (2)$$

since both the spectrum shape and the polarization are dependent on these two parameters only (the unexplained symbols are defined by Konopinski[11]). For a quantitative discussion of the polarization measurement one has to determine ξ_1 and η_1 by fitting the experimental (2) spectrum shape factor to the theoretical one and to calculate the polarization using these values. To get a rough idea we can use the approximation $\alpha Z \ll 1$, $p\rho \ll 1$, which really is not well satisfied by RaE. Then the polarization can be written conveniently as

$$P = -\frac{v}{c}\,\frac{X(W,\xi_1,\eta_1) - Y(W,\xi_1,\eta_1)}{X(W,\xi_1,\eta_1) - (\frac{v}{c})^2 Y(W,\xi_1,\eta_1)}. \qquad (3)$$

(W is the electron energy, the functions X and Y can be obtained with the help of reference 7.) It can be seen that a deviation of $-P/(v/c)$ from 1 must occur for $v/c \neq 1$.

A more detailed description of this experiment can be found in reference 5.

Thanks are due to Professor O. Haxel for his interest in this work.

† Supported in part by the Deutsche Forschungsgemeinschaft.

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DETERMINATION OF THE PARITIES OF STRANGE PARTICLES FROM DISPERSION RELATIONS*

Saul Barshay†
Brookhaven National Laboratory,
Upton, New York
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Several authors [1-3] have recently suggested that the parities of the K-Λ and K-Σ systems relative to the nucleon might be determined

from the forward angle dispersion relations applied to the scattering of K^+ and K^- mesons by nucleons. The dispersion relations together with detailed measurements on the pion-nucleon system have provided an important test of the principle of microscopic causality in field theory. As applied to the K meson-nucleon system the relations will not, as yet, provide a theory capable of much predictive power as to the dynamics, but if we can use experimental results together with them to determine the relative parities of the strange particles an important step will have been taken toward the construction of a possible dynamical theory.

Igi[4] has recently analyzed the present data on K meson-nucleon interactions on the basis of the following dispersion relation[5,6]:

$$2F=(\omega/K^2)\{D_+(\omega)-\tfrac{1}{2}(1+\omega/m_K)D_+(m_K)-\tfrac{1}{2}(1-\omega/m_K)D_-(m_K)\} - \frac{\omega}{4\pi^2}\int_{m_K}^{\infty}\frac{d\omega'}{K'}\left\{\frac{\sigma_+(\omega')}{\omega'-\omega}+\frac{\sigma_-(\omega)}{\omega'+\omega}\right\} - \frac{\omega}{\pi}\int_{\omega_0}^{m_K}d\omega'\,\frac{A_-(\omega')}{K'^2(\omega'+\omega)}\,.$$

Here the $D_\pm$ are forward scattering amplitudes for K^+ and K^-, respectively, of total laboratory energy ω on protons, the $\sigma_\pm$ are total cross sections, m_K is the K meson mass, ω_0 is $\sim 0.4 m_K$, and $A_-(\omega)$ is the absorptive part of the K^--p scattering amplitude. Let us choose the convention that the Σ- nucleon relative parity is positive. Then on the assumption that the Σ-Λ relative parity is positive, Igi[4] uses the following values for $2F$:

$$\text{(a) } 2F=-g_1^2\left(\frac{1}{2M}\,\frac{1}{M_1-M-\omega_1}\right)-g_2^2\left(\frac{1}{2M}\,\frac{1}{M_2-M-\omega_2}\right)$$
$$=-(1.01\,g_1^2+0.64\,g_2^2)\,m_K^{-2} \text{ for scalar } K,$$
$$\text{(b) } 2F=g_1^2\left(\frac{1}{2M}\,\frac{1}{M_1+M+\omega_1}\right)+g_2^2\left(\frac{1}{2M}\,\frac{1}{M_2+M+\omega_2}\right)$$
$$=(0.062\,g_1^2+0.060\,g_2^2)\,m_K^{-2} \text{ for pseudoscalar } K. \quad (2)$$

Here M_1, M_2 and M are the Λ, Σ, and nucleon masses, respectively, and g_1 and g_2 are the renormalized couplings of Λ and Σ particles, respectively, to nucleons and K mesons; $\omega_1 \sim 0.107 \times m_K$. and $\omega_2 = 0.268\, m_K$. We now want to drop the assumption of positive Σ-Λ relative parity. If this relative parity is negative then Eq. (2) is changed to read

$$\text{(a) } 2F = 0.062\,g_1^2-0.64\,g_2^2 \text{ for scalar } K,$$
$$\text{(b) } 2F = 1.01\,g_1^2 + 0.060\,g_2^2 \text{ for pseudoscalar } K. \quad (3)$$

Now we note immediately that if g_1^2 and g_2^2 do not differ by an order of magnitude or more (which is certainly quite probable[7]) Eqs. (3a) and (3b) both give a negative F. As noted previously,[1-3] Eq. (2a) gives a negative F and Eq. (2b) a positive F. As noted by Salam and Matthews,[3] the relative K-Σ parity might be determined rather unambiguously from the dispersion relation for K^+-n scattering analogous to Eq. (1). Then only the Σ intermediate state contributes to the corresponding F, which if negative implies a scalar K, and if positive implies a pseudoscalar K. To what extent does the sign of F in the K^+-p dispersion relation determine the second unknown relative parity? Igi's analysis[4] shows that if the effective K^+-p potential is repulsive and if the K^+-p cross section is essentially isotropic and energy independent from 0 to ~200 Mev kinetic energy, then F is negative for either a repulsive or an attractive K^--p effective potential.[8] From Eqs. (2) and (3) we see that only if the relative K-Σ parity is negative does a negative F determine the relative Σ-Λ parity (to be negative). If the relative K-Σ parity is positive then a negative F is compatible with either relative Σ-Λ parity. Stated another way, if we had by some method determined the relative Σ-Λ parity, then only if this turned out to be positive would a negative F determine the relative K-Σ parity (to be positive). On the other hand we note if F were to turn out to be positive, Eq. (3b) would imply pseudoscalar K mesons and positive relative Σ-Λ parity.

In the event that a future analysis of K^+-n scattering should indicate a positive K-Σ relative parity[9] then the relative Σ-Λ parity may still have to be determined by some experiment such as the angular variation of the polarization of the Λ from the reaction $\Sigma^- + p \to \Lambda + n$, as suggested by Pais and Treiman,[10] or by the internal conversion coefficient in Σ^0 decay as calculated by Feinberg[11] and by Feldman and Fulton.[12] The latter effect is, unfortunately, almost insensi-

tive to the relative Σ-Λ parity. Both experiments are quite difficult from the standpoint of the accumulation of statistics alone. One might still try to appeal to a dispersion relation at this point. Consider elastic π^0-Λ scattering. A dispersion relation for this process with no subtraction may be written

$$D_0(\omega)=\alpha+(1/2\pi^2)\int_0^\infty K'^2 dK' \sigma_0(K')/(K'^2-K^2), \quad (4)$$

where D_0 is the forward scattering amplitude and σ_0 is the total cross section; α is given by $\alpha=-2b^2\omega_0/(\omega^2-\omega_0^2)<0$ for negative relative Σ-Λ parity and $\alpha=[2b^2/(M_1+M_2)^2][\omega_0(\mu^2-\omega_0^2)/(\omega^2-\omega_0^2)]>0$ for positive relative Σ-Λ parity. Here μ is the pion mass, $\omega_0=(M_2^2-M_1^2-\mu^2)/2M_1\sim 70$ Mev, and b is the renormalized coupling of a pion to a Σ and a Λ. Only the Σ^0 intermediate state contributes to α. We cannot scatter pions on hyperons directly, but two potentially useful reactions that produce high-energy and low-energy pion-hyperon states, respectively, are $K^- + p \rightarrow \Lambda + \pi^0$ and $K^- + p \rightarrow \Lambda + \pi^0 + \pi^0$. It may be that future experiments on these reactions together with a theoretical analysis could determine the sign of the low-energy scattering length in the Λ-π^0 system.

Finally, we note that Eq. (3), as well as Eq. (2),[2-4] is consistent with renormalized K-meson couplings with $g^2 \sim 1-4$.

* Work performed under the auspices of the U. S. Atomic Energy Commission.

† A National Science Foundation Postdoctoral Fellow.

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ANTI-LAMBDA HYPERON*†

D. J. Prowse‡
University of California at Los Angeles, California
and
M. Baldo-Ceolin
University of Padova, Padova, Italy
(Received August 7, 1958)

An event has recently been found in an emulsion stack exposed to a 4.6 ± 0.3 Bev π^--meson beam at the Berkeley Bevatron which is interpreted as the decay of an anti-lambda hyperon, $\bar{\Lambda}^0$. The threshold for production with a free nucleon is 4.73 Bev and extends down to ~ 4.3 Bev with a bound nucleon, the production reaction being of the type

$$\pi^- + p \rightarrow \Lambda^0 + \bar{\Lambda}^0 + n.$$

The side of the stack was exposed to a beam of 10^6 π mesons per sq cm, the intensity falling off by a factor of 10^3 in a distance of 8 in. across the emulsion sheets. The emulsions were scanned for stopping π^+ mesons and these were followed back to their origins if it appeared that they were produced in the region of high π-meson flux. Most tracks which were successfully traced back came from primary π^--meson stars but 3 traced back to decaying K^+ mesons and one to a "V" type event. The event is shown in Fig. 1. The other branch of the "V" event was traced after 3 cm range into a large star from which 3 shower particles were emitted. The visible energy in this star is 783 Mev including the rest masses of the 3 shower particles which were all identified as π mesons; one was arrested and decayed into a μ meson, the other two were identified by the Δg^* vs ΔR and by the $\bar{\alpha}$ vs g^* methods.

The particle connecting the "V" event with the large star has been shown to have been traveling into the star rather than the reverse by ionization measurements on the track. These are summarized in Table I. Every useful grain in the track has been counted and the plateau ionization in each plate has been determined by measuring the beam π mesons. The grain den-

Table I. Details of ionization measurements on the antiproton track.

Plate	Total corrected range	Range counted	No. grains	Minimum	g^*	Effective range from $\bar{\Lambda}^0$-decay
1	2.78 mm	2.27 mm	718	15.00	2.104	1.13 mm
2	5.88	4.72	1572	15.51	2.150	5.72
3	6.59	5.11	1780	15.62	2.230	11.98
4	6.78	5.42	1890	15.39	2.260	18.67
5	5.09	4.56	1601	15.19	2.316	24.84

sities have then been normalized to the minimum ionization value via the plateau[1] in each plate. It is clear that the track increases in ionization towards the star. The star has therefore been identified as an antiproton interaction in flight. The best value of the antiproton energy at the decay point is 230_{-7}^{+22} Mev. The opening angle of the V is $64 \pm 1^\circ$ and the π^+-meson energy is 32 Mev from a range of 1.70 cm in emulsion.[2] From these values the Q value in the decay has been calculated to be $35_{-0.9}^{+2.6}$ Mev, this is in excellent agreement with that for the normal Λ^0 hyperon: 37.45 Mev.[3]

One cannot of course completely rule out the possibility that the decay event is really the interaction of an antineutron charge-exchanging into an antiproton and also creating a π^+ meson, but the visible energy in the star, coupled with the direction of the connecting track and the apparent Q value in the decay, makes the $\bar{\Lambda}^0$ hyperon the only reasonable interpretation.

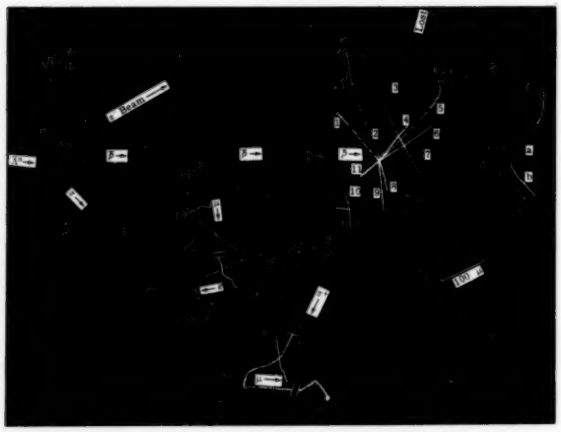

FIG. 1. Drawing of the event. The $\bar{\Lambda}^0$ hyperon enters from the left in the direction indicated. The antiproton track is shown and is almost horizontal. The star apparently caused by the interaction of the antiproton in flight has 11 prongs, all but one of which have been arrested in the emulsion. Of the three shower tracks, one interacts in flight and a second decays at rest. The π^+-meson track resulting from the decay of the antihyperon is shown decaying into a μ meson and thence to an electron.

We would like to thank Dr. E. J. Lofgren and his colleagues at the Bevatron for making the exposure possible. We are grateful to Professor D. H. Wilkinson and Dr. W. Chupp who carried out the exposure. We are indebted to Professor C. F. Powell, Professor N. Dallaporta, Dr. D. Evans, Dr. D. H. Stork and Professor W. F. Fry for many very useful discussions.

*A very preliminary report of this event was made at the Washington Meeting of the American Physical Society in May, 1958 [M. Baldo-Ceolin and D. J. Prowse, Bull. Am. Phys. Soc. Ser. II, 3, 163 (1958)]. A more detailed account than is given in this letter will be published elsewhere.

†Partially supported by the United States Atomic Energy Commission.

‡On leave from the University of Bristol, Bristol, England.

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BERYLLIUM-7 EXTREME HIGH-TEMPERATURE EXPERIMENT

Carl A. Accardo
United States Signal Research and Development Laboratory, Fort Monmouth, New Jersey.
(Received August 4, 1958)

The recent attainment of extreme high temperatures of milliseconds duration suggests a rather interesting experiment.[1] It is known that Be^7

decays by a 100% K-capture process with a 52-day half-life. The decay proceeds in two branches; 88% of the transitions go directly to the ground state of lithium, and the remaining 12% to an excited state of lithium followed by a 0.48-Mev gamma ray. The Be^7 photopeak is easily detectable by using a conventional NaI crystal spectrometer in view of the large photoelectric-to-Compton cross section ratio at this energy.

Since the fourth ionization potential in Be is 216 ev, one would expect that the introduction of the isotope into a high-temperature discharge (2.5×10^6 to 10^7 degrees Kelvin; 216 ev to 860 ev) would result in complete ionization of the Be. By monitoring the 0.48-Mev line with the proper discrimination and shielding against the softer x-rays, one should observe the stripping of the K electrons for the duration of the high temperature by measuring the drop in the number of counts in the photopeak. Assuming high-temperature times of millisecond duration, the experiment could be carried out with a millicurie source if properly introduced into the discharge tube. The experiment would be of interest as a possible method of high-temperature thermometry in the million-degree range and also provide information regarding electron-nucleus interaction in K-capture processes.

[1] P. C. Thonemann *et al.*, Nature 181, 217 (1958).

TRAPPED ALBEDO THEORY OF THE RADIATION BELT

S. F. Singer
Department of Physics,
University of Maryland,
College Park, Maryland.
(Received August 12, 1958)

We describe here a theory for the radiation belt of the earth which gives the energy spectrum, angular distribution, and intensity of the radiation with distance from the earth. Our results are not compatible with data reported by Explorer satellites I and III (where the observed intensity did not rise appreciably below 1000 km)[1] but are compatible with preliminary reports from Explorer IV.[2]

The earth's dipole field provides an excellent trapping field not only for the magnetic storm-producing solar corpuscular radiation but also for cosmic-ray albedo particles.[3] Because of the long lifetimes which are possible we must investigate all likely injection mechanisms. From (12*) and (13*) it is evident that, e.g., high-energy electrons could stay trapped for long times. However, we have found no adequate production mechanism.[4]

On the other hand, the injection of protons arising from the decay of upward-projected neutrons gives an adequate value of Q. Because of the long neutron lifetime, we can have copious neutron production in stars in the dense lower atmosphere with protons released in regions of very low density. The release rate is not very altitude sensitive, leading to constant Q referred to earlier,[3] and therefore to the possibility of using the radiation belt as a tool for study of the exosphere.

The situation has been considered in detail for the equator up to ~1 earth radius (R_E=4000 miles), since experiments, e.g., a rocket measurement of flux vs altitude, could be analyzed most simply.

Within experimental uncertainties, the primary flux J is about $10^{-1}\text{cm}^{-2}\ \text{sec}^{-1}$. Each primary is assumed to produce 7 stars[5]; we estimate that ~2.5 are effective, i.e., sufficiently high in atmosphere to allow neutrons to escape. The neutron spectrum for the average star in nitrogen is given as [6,4] (derived from Bristol star data)

$$i_n(E)=8E^{-1.6}\ \text{per Mev or}\ i_n(\beta)=0.12\ \beta^{-2.6}. \tag{1}$$

Thus the spectrum of neutrons emitted upwards per cm^2 per sec is

$$J_n(\beta)=0.03\ \beta^{-2.6}\equiv C_1\beta^{-\gamma}. \tag{2}$$

Since the effect of absorption and decay on the neutron flux can be neglected, we are justified in assuming a constant "brightness" with altitude, at least for a fraction of R_E. Thus the source function of decay protons $Q(\beta)$ can be assumed uniform over moderate altitudes; it is also reasonably isotropic:

$$Q(\beta)=(\Lambda/c)C_1\beta^{-(\gamma+1)}=C_2\beta^{-(\gamma+1)}, \tag{3}$$

where $\Lambda=9\times10^{-4}\ \text{sec}^{-1}$.

We next investigate the fate of the protons: (i) Coulomb scattering which leads to diffusion in α and leakage from the trap[3]; and compare it to (ii) energy loss by collisions. We will find the latter to be more important for protons with $\beta<1$.

For the nonrelativistic case we find the mean

square scattering[7] in pitch angle

$$(\Delta\alpha^2)/t=\tfrac{1}{2}\{16\pi r_e^2 N(Z^2/A)(mc/\beta p)^2 \times \ln[196Z^{-\frac{1}{3}}(Z/A)^{\frac{1}{3}}]\}\, cd\cdot\beta = 1.2\,(Z^2/A)(m/M)^2 \ln[\,]\, cd\cdot\beta^{-3} \text{ per second.} \quad (4)$$

$(\Delta\alpha^2)/t$ is defined as the diffusion coefficient D [see Eq. (4*)]. The relativistic approximation to D [Eq. (13*)] may be compared with the nonrelativistic set of values

$$D_{in}=C_{in}\, cd\cdot\beta^{-3};\quad D_{out}=C_{out}\, cd\cdot\beta^{-3}, \quad (5)$$

where $C_{in}=5.6\times10^{-6}$ and $C_{out}=1.9\times10^{-6}$.

We next evaluate the velocity loss per second for nonrelativistic protons:

$$-d\beta/dt = (cd/Mc^2)(-dE/dx),$$

where $(-dE/dx)$ is the collision loss[7] per g-cm^{-2}. we find

$$-d\beta/dt \cong -2.6\times10^{-3}(Z/A)\, cd\cdot\beta^{-2}\equiv -C_3\beta^{-2} \quad (6)$$

from which

$$\beta^3=\beta_0^3-0.004\, cd\cdot t\,. \quad (7)$$

We see that as β decreases in accordance with (7), D will increase [see Eq. (5)]. The time for β to decrease over the maximum region of interest, from $\beta\cong0.75$ (500 Mev) to $\beta=0.2$ (20 Mev) is $\sim100(cd)^{-1}\beta_0^3$. At 1000 km altitude $t\sim3\times10^{-10}$ sec. On the other hand, the mean lifetime T is given by (12*):

$$T\sim2(\sim0.1)/D\sim4\times10^4(cd)^{-1}\beta^3\,.$$

Since $T\gg t$, we will to a first approximation neglect leakage due to diffusion in pitch angle and consider only momentum loss by collisions.

To derive the energy spectrum of the trapped protons we consider an equation of continuity in "β space,"

$$\partial N/\partial t+\operatorname{div} j(\beta)=Q(\beta)\,, \quad (8)$$

where we define the flow $j(\beta)=N(\beta)d\beta/dt$;

$$\operatorname{div} j=\frac{dj}{d\beta}=\frac{d}{d\beta}\,[c_\nu\beta^{-\nu}(-C_3\beta^{-2})] = \frac{C_\nu C_3}{\nu+2}\,\beta^{-(\nu+3)} \quad (9)$$

To evaluate $N=C_\nu\beta^{-\nu}$ we use the steady state condition $\partial N/\partial t=0$. Then

$$N(\beta)=C_2C_3^{-1}\gamma\beta^{-\gamma+2}\equiv C_4\beta^{-\gamma+2} \quad (10)$$

The omnidirectional flux is obtained by integrating over the allowed solid angle, as defined by α_C:

$$J(\beta_2,\beta_1)=\int_{\beta_1}^{\beta_2}N(\beta)\beta c d\beta\int d\Omega = C_4c(4-\gamma)^{-1}\,(\beta_2^{4-\gamma}-\beta_1^{4-\gamma})\times 4\pi\cos\alpha_C. \quad (11)$$

Substituting from (3), (2), (6), (10) we have $\gamma=2.6$, $C_1=0.03$, $C_2=9\times10^{-16}$, $C_3=1.56\times10^8 d$, $C_4=1.5\times10^{-23}d^{-1}$; therefore

$$J(\beta_2,\beta_1)=(C_5\cos\alpha_C)f(\beta_2,\beta_1) =4\times10^{-12}(d^{-1}\cos\alpha_C)\,f(\beta_2,\beta_1). \quad (12)$$

Our final result for the equatorial flux is contained in Eq. (12). The energy spectrum is seen to be extremely flat. The function $f(\beta_1,\beta_2)$ is plotted in Fig. 1 for three values of β_2. We also show the amount of lead shielding necessary to reduce the radiation intensity. A reduction by a factor 2 requires 10 g-cm^{-2} (or 1 cm) of Pb, whereas a 10-fold reduction requires about 4 cm of Pb. The energy spectrum will fall off more

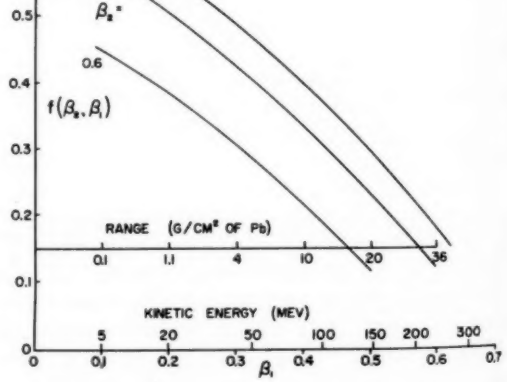

FIG. 1. $f(\beta_2,\beta_1)\equiv\beta_2^{1.4}-\beta_1^{1.4}$ is proportional to the number of protons between β_2 and β_1. For an ordinary counter, $\beta_1\sim0.1$ (or 5 Mev). With an absorber around the detector, β_1 can be raised (as shown), thereby reducing the flux.

steeply beyond 400 Mev for several reasons: (i) the spectrum (1) steepens; (ii) the proton range is about one mean free path for nuclear collisions; (iii) the adiabatic condition (2*) breaks down.[3]

The altitude dependence is given by $\sim 1/d$ (to a first approximation). For example, at 1000 km altitude $d \sim 10^{-19}$ g/cm^3. With $f(\beta_2, \beta_1) \sim 0.5$ and[8] $\cos \alpha_C \sim 0.34$, we get for the omidirectional flux

$$J(1000\text{km}) \sim 7 \times 10^6 \text{ particles/cm}^2\text{-sec.} \quad (13)$$

(Any discrepancy with a measured flux value may therefore indicate that the density at 1000 km is different from the value estimated above.)

In order to obtain a more accurate theory for the variation of J with altitude h, we discuss some of the approximations entering the derivation of Eq. (12). We had implicity assumed d = constant in the region in which the particle moves. This is certainly not the case for small α_0 since there the particle can move far along the line of force, and therefore deep into the denser atmosphere. Therefore we introduce an approximate value of $\cos\alpha_C$ in the manner shown.[8] But there is no sharp cutoff. But even particles with $\alpha_0 \sim \pi/2$ (which remain in the equatorial plane) will traverse a varying d, if their radius of curvature $\rho \gtrsim H_0$, i.e., for energy $\gtrsim 50$ Mev ($\beta \gtrsim 0.3$).

The situation can now be summarized as follows: the decay protons are released, to a good approximation, isotropically and uniformly with altitude. Their energies range up to several hundred Mev, corresponding to ρ up to ~ 100 km. Below the exosphere the lifetime is very short because ρ is close to h and the path length becomes large.[9] Hence the intensity of trapped radiation can only rise in the exosphere (where also H_0 becomes $\gtrsim \rho$). If we take α_C independent of altitude h (i.e., assuming a constant scale height H_0), then

$$J \propto \exp[(h - 400)/H_0]. \quad (14)$$

The scale height is given by $kT/\mu Mg$. T is constant (by Liouville's theorem) and $\sim 1000^\circ$K, but not well known. μ decreases from about 20 at 400 km (corresponding to a mixture of N_2, N, and O) to 1 (corresponding to H atoms) above the transition altitude. Thus H_0 increases from 50 km up to 1000 km. Correspondingly, we would expect a moderate rise initially (more rapid except for the protons with $\rho \gtrsim H_0$), which would be followed by a slower rise when hydrogen predominates. (However this rate of rise is somewhat enhanced because $\cos\alpha_C$ increases with H_0.) For example, at around 500 km where H_0 may be ~ 200 km, we expect the flux to increase by ~65% in a 100-km interval.

It is interesting to speculate on the variation of J out to several earth radii. Beyond 1000 km, d can decrease by not more than 2 or 3 powers of 10; the gas becomes more ionized so that the ordinary scale height concept loses significance. The factor $\cos\alpha_C$ can increase by only a factor of 3. On the other hand Q must fall off as R^{-2}. We therefore anticipate a broad maximum in the equatorial plane at 1-2 R_E followed by a decrease as R^{-2}, and perhaps faster. The forthcoming "moon shots" should be able to check this prediction.

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1 Van Allen, Ludwig, Ray, and McIlwain, preliminary experimental results from U.S. International Geophysical Year Satellite 1958 alpha and gamma presented at the National Academy of Sciences, May 1, 1958 (unpublished).

2 New York Times, August 10, 1958.

3 S. F. Singer, accompanying Letter [Phys. Rev. Lett. 1, 171 (1958)]. The same notation is used. The asterisk refers to equations and figures in that Letter.

4 Griem, Singer, and Wentworth (to be published).

5 G. Puppi and N. Dallaporta, in Progress in Cosmic-Ray Physics (Interscience Publishers, New York, 1952), Vol. 1, p. 379.

6 R. C. Wentworth and S. F. Signer, Phys. Rev. 98, 1546 (A) (1955).

7 B. Rossi, High-Energy Particles (Prentice-Hall, New York, 1952).

8 We calculate $\cos\alpha_C$ by setting a lowest level (along the line of force) to which particles may penetrate and survive for long times. (A smilar situation holds in the case of pure scattering; see Fig. 1*.) If we set $\Delta h = H_0$ the scale height (~ 200 km), we find $\cos^2\lambda_* = 1 - \Delta h/(R_E + h)$; and from (3*), $\cos\alpha_C = 0.34$. This corresponds to $\alpha_C \sim 70^\circ$.

9 H. Griem and S. F. Singer, Phys. Rev. 99, 608 (A) (1955). See also Fig. 13c of S. F. Singer, in Progress in Elementary Particle and Cosmic Ray Physics (Interscience Publishers, New York, 1958), Vol. 4, p. 263.

ERRATA

FURTHER OBSERVATIONS ON NATURE OF THE CURRENT REDUCTION IN THE PRIMARY COSMIC RAY INTENSITY. Martin A. Pomerantz and Som Prakash Agarwal [Phys. Rev. Lett. 1, 107 (1958)].

In the last column of Table I, the second entry

from the bottom should read 16.0 ± 0.5 instead of 18.4 ± 1.3.

CORRELATION AND THE ISOTOPE EFFECT FOR DIFFUSION IN CRYSTALLINE SOLIDS. A. H. Schoen [Phys. Rev. Lett. 1, 138 (1958)].

The word "interstice" should be replaced wherever it occurs by "interstitialcy". This change was made without the author's knowledge.

On page 139, in the third line above Eq. (7), the last word should be "or" rather than "of."

ABSTRACTS

In this section are printed the abstracts of Articles that have been forwarded to The American Institute of Physics for publication in THE PHYSICAL REVIEW. In quoting information obtained from this section before the appearance of the corresponding Article, reference should be made to "Physical Review (to be published)" rather than to this Journal.

STRUCTURE OF A SHOCK WAVE IN FULLY IONIZED HYDROGEN. D. A. Tidman, The Enrico Fermi Institute for Nuclear Studies, The University of Chicago, Chicago, Illinois (Received March 24, 1958).

The Fokker-Planck equations are used to examine the structure of a shock wave in fully ionized hydrogen. This is done by assuming a bimodal Maxwellian distribution for the protons in the interior of the shock and noting that the electrons are in thermal equilibrium with themselves but not necessarily with the protons. The method is essentially an extension of that used by Mott-Smith in his analysis of the Boltzmann equation for a shock wave in a gas of neutral atoms.

LENGTH EFFECT IN THE HEAT TRANSPORT IN HELIUM II. H. Forstat, Michigan State University, East Lansing, Michigan (Received May 20, 1958; revised manuscript received June 13, 1958).

A length effect in the heat transport in Helium II in small channels has been observed. Measurements were made using a column of packed jeweller's rouge for lengths of 3.179, 5.166, and 8.156 cm in the temperature range 1.70-2.17°K. The channel size was approximately 0.1 micron. The results verified the linear dependence of the heat current density on temperature difference and showed a decrease in the heat transport by approximately a factor of 2 as the rouge length was increased from 3.179 to 8.156 cm.

MAJORANA FORMULA. Alvin Meckler, Division of Physical Sciences, National Security Agency, Fort Meade, Maryland (Received March 14, 1958).

The Majorana formula for the transition probability of a general spin is here rederived and brought into a more compact form. The method is based on the use of the projection operators which characterize the spin vector as definitely quantized along a specified unit vector. If the Hamiltonian contains only terms linear in the spin, it is possible to define a moving instantaneous axis along which the moving spin vector maintains its quantization. The Majorana problem is reduced to the calculation of the joint probability of quantization along two different axes.

The usual expression of the Majorana formula contains a factor which is the transition probability of a spin-$\frac{1}{2}$ vector. It is shown here how this factor is related to the angle through which the general spin is turned.

SUPRATHERMAL PARTICLES. E. N. Parker and D. A. Tidman, The Enrico Fermi Institute for Nuclear Studies and the Department of Physics, The University of Chicago, Chicago, Illinois (Received May 1, 1958).

The production of suprathermal particles in agitated plasmas bearing magnetic fields is discussed. This is done by setting up a Fokker-Planck equation to describe the phenomenon. The role of such particles in astrophysics and in the problem of the production of thermonuclear power is considered.

BOMBARDMENT OF CADMIUM SULFIDE CRYSTALS WITH 30- TO 60-kev ELECTRONS. C. E. Bleil, D. D. Snyder, and Y. T. Sihvonen, General Motors Research Staff, Detroit, Michigan (Received May 9, 1958).

The dependence of induced conductivity in CdS crystals on the rate of arrival and the energy of impinging electrons is reported. These results lead to a qualitative picture of the conduction, excitation, and recombination phenomena in CdS which is satisfied by the simple model of a sulfur vacancy, and an analytical expression which involves the mobility of carriers, their effective masses, the number of ground states in the forbidden gap (for pure but not perfect crystals), and their positions as a function of temperature. The red and green luminescence observed under irradiation by electrons is qualitatively explained. The equation for the induced conductivity is derived from the expression for two-carrier conductivity and the definitions of the steady state Fermi levels for holes and elec-

trons. An empirical fit to the observed data yields $\sigma = e\mu_n N_c \exp[(kT\alpha I - E_{fn}^0)/kT]$ for $I < I_S$ and $\sigma = e(\mu_n + \mu_p) N_v \exp\{[kT\beta(I - I_S) - E_{fp}^0]/kT\}$ for $I > I_S$. The data reported herein were obtained using monoenergetic bombarding electrons in the range of 30 to 60 kev. An interesting field effect is reported but no clear interpretation is available to the authors. Corroborating data employing ultraviolet irradiation is presented.

THEORY OF THE PIEZOELECTRIC EFFECT IN THE ZINC BLENDE STRUCTURE. Joseph L. Birman, Research Laboratories, Sylvania Electric Products, Incorporated, Bayside, New York (Received March 18, 1958).

A theory of the piezoelectric effect in the zinc blende structure is developed on the basis of two assumptions: (1) the Born theory of a homogeneous lattice deformation, (2) that each chemical bond is characterized by a vector bond dipole moment. The theory leads to an equation relating the macroscopic piezoelectric constant to the static and dynamic effective charges (properties of the bond moment), and a lattice parameter γ_{14} which relates the internal strain (basis deformation) to the external strain. An x-ray method for measuring γ_{14} is discussed; it utilizes the failure of Friedel's Rule in the anomalous dispersion region. An order of magnitude check of the theory is made on ZnS. Some remarks are made about the piezoelectric effect in wurtzite based on the same model.

ELECTROLUMINESCENCE OF ZnS SINGLE CRYSTALS WITH CATHODE BARRIERS. D. R. Frankl, Central Research Laboratories, Sylvania Electric Products, Incorporated, Bayside, New York (Received April 29, 1958).

Certain ZnS crystals show electroluminescence predominantly near the cathode. The major emission from Cu-activated crystals occurs as a burst of light when the exciting voltage is suddenly removed, as electrons flowing back toward the cathode to neutralize the barrier charge recombine with ionized luminescent centers. This burst may be quenched by a voltage in the initial direction or enhanced by a voltage in the opposite direction. There is also a small steady component of emission due either to capture of electrons that entered the crystal from the cathode or to excitation of centers to discrete levels. In Mn-activated crystals, only the latter has been observed.

These results, as well as the emission peaks obtained from copper-activated ZnS under sinusoidal voltage excitation, are explained in terms of ionization of luminescent centers in a barrier region. The "primary" peak, which predominates in powder samples and under low voltages in single crystals with metal contacts, is due to recombination of centers left over from the previous half-cycle. A capture cross section of about 10^{-12} cm^2 is deduced. The peak at about 150 to 170 phase degrees is due to recombination of ionized centers left in the neutral part of the crystal as the barrier width diminishes. A barrier space charge density of (very roughly) 10^{16} e/cm^3 is deduced. The in-phase peak is analogous to dc electroluminescence.

SUPERCONDUCTING TO NORMAL PHASE TRANSITION IN TANTALUM. William B. Ittner, International Business Machines Corporation, Poughkeepsie, New York (Received March 26, 1958).

Measurements were made of the rate at which superconducting phase collapses radially in a hollow cylindrical tantalum specimen following the sudden application of a longitudinal magnetic field greater than the critical field. The measured transition rates confirm the hypothesis that the propagation is controlled by electromagnetic damping associated with eddy currents generated in the normal phase. The results, moreover, may be interpreted on the basis of a theoretical treatment of the problem first published by Pippard, provided that suitable modifications are incorporated to include the thermal effects which accompany the transition.

F^{19} HYPERFINE INTERACTION IN THE PARAMAGNETIC RESONANCE SPECTRUM OF Mn^{++} IONS IN ZnF_2. A. Mukherji and T. P. Das, Saha Institute of Nuclear Physics, Calcutta, India (Received September 30, 1957; revised manuscript received June 10, 1958).

The observation of a predominantly isotropic fluorine hyperfine interaction in the paramagnetic resonance of Mn^{++} ion dissolved in ZnF_2 crystal has been interpreted in terms of the contact interaction of the unpaired d electrons of Mn^{++} ion at the fluorine nucleus. The distortion

produced in the free-ion orbitals of Mn^{++} $3d$ electrons due to the electrons of the F^- ion has been considered. The calculated value of the isotropic hyperfine interaction constant is 10 oersteds, which is about half the experimental value of 19 oersteds. The quadrupole interaction of the Mn^{++} d electrons with the crystalline field gradient calculated on the point-charge model is found to produce a very small change in the theoretical value. The principal causes which contribute to the quantitative disagreement between the theoretical and the experimental values are also discussed.

V-TYPE CENTER RESONANCE OF NEUTRON-IRRADIATED LiF AT ROOM TEMPERATURE. Y. W. Kim, R. Kaplan, and P. J. Bray, Department of Physics, Brown University, Providence, Rhode Island (Received May 1, 1958).

A V-type center similar to one described by Känzig and co-workers has been found in LiF single crystals neutron-irradiated at room temperature. Analysis of the s and p admixture in the electron wave function, using the Hamiltonian of Woodruff and Känzig, is made from the angular dependence of the electron-spin resonance spectrum.

EFFECT OF MONOENERGETIC FAST NEUTRONS ON N-TYPE GERMANIUM. S. L. Ruby, F. D. Schupp, and E. D. Wolley, Materials Engineering Departments, Westinghouse Electric Corporation, East Pittsburgh, Pennsylvania (Received May 12, 1958).

The density of vacancy-interstitial pairs produced per unit of fast neutron flux on n-type germanium has been measured experimentally using monoenergetic neutrons. The changes in resistivity and Hall coefficient have been determined as a function of neutron flux and energy for n-type germanium having impurity concentrations of about $2\times10^{12}cm^{-3}$. Both the irradiation and electrical measurements have been made at about 77°K. Neutron energies in the experiment were 4.8, 3.2, and 1.7 Mev from the $H^2(d,n)He^3$ reaction and the observed changes in resistivity and Hall coefficients indicate charge carrier removal rates per unit of neutron flux, $-dn/d\phi$, of 12, 12, and 21 respectively. Not only do the results appear much lower than the value of 160 calculated by current models, but are less than would be estimated by extrapolating the results obtained with reactor neutron irradiations.

RECOMBINATION PROPERTIES OF GOLD IN SILICON. G. Bemski, Bell Telephone Laboratories, Incorporated, Murray Hill, New Jersey (Received May 21, 1958).

The presence of gold atoms in the silicon lattice decreases the lifetime of excess electrons and holes in p- and n-type material. The capture of electrons in p-type silicon occurs through the gold donor level with a capture cross section, σ_{no}, of $3.5\times10^{-15}cm^2$ (at 300°K). This capture cross section varies as $T^{-2.5}$ between 200° and 500°K. In n-type silicon, the electron capture cross section $\sigma_{no}(300°K) = 5\times10^{-16}cm^2$ and is temperature independent; the hole capture cross section $\sigma_{po}(300°K) = 1\times10^{-15}cm^2$ and varies as T^{-4}. The capture in this case occurs through the gold acceptor level.

PROPAGATION OF A MAGNETIC FIELD INTO A SUPERCONDUCTOR. Joseph B. Keller, Institute of Mathematical Sciences, New York University, New York, New York (Received May 28, 1958).

The propagation of a magnetic field into a superconducting wire of circular cross section is analyzed theoretically. Upper and lower bounds are obtained on t_0, the time required for the field to reach the axis, and on $R(t)$, the inner radius of the normally conducting region. The lower bound is just the approximate result obtained by Sixtus and Tonks, Pippard, and Lipshitz. It has been verified experimentally by Farber for applied fields less than 1.04 times the critical field. It is believed that for large applied fields the actual results should be closer to the upper bounds.

FERRIMAGNETIC RESONANCE IN YTTRIUM IRON GARNET AT LIQUID HELIUM TEMPERATURES. J. F. Dillon, Jr., Bell Telephone Laboratories, Murray Hill, New Jersey (Received May 20, 1958).

Ferrimagnetic resonance experiments have been performed on single-crystal spheres of yttrium iron garnet $[Y_3Fe_2(FeO_4)_3]$ at 24 kMc/sec. At 4.2°, the line width measured with the steady field along [100] is 15 oe, as compared with 30

oe at 295° and 400 oe at 65°K. At liquid helium temperatures, the crystal is apparently cubic. There is a strong anisotropy in line width. Except for the case of the field along [100], the line width varies strongly with temperature in this temperature region. A large magnetocrystalline anisotropy is indicated by the variation in field for resonance with direction. This ranges over 3000 oe. The maxima and minima in this field indicate that terms of higher order than the K_1 and K_2 terms used hitherto will be necessary for the description of the anisotropy energy surface. Thus, at 4.2°K, apparently [100], [110], and [111] are all hard directions, and the easy directions are exactly or close to [112], [221], and [123].

FURTHER STUDIES ON SOME INFRARED STIMULABLE PHOSPHORS: QUENCHING, STIMULATION, AND EXHAUSTION STUDIES. Seymour P. Keller and George D. Pettit, Research Laboratory, International Business Machines Corporation, Poughkeepsie, New York (Received May 8, 1958).

As a continuation of earlier work on the two phosphor systems SrS:Ce, Sm and SrS:Eu, Sm, studies of the quenchings and the enhancements of fluorescent emission by visible and infrared radiation have been made. Two stimulation bands have been measured and the relative efficiencies of exhaustion and of stimulation of visible emissions by the two bands have been studied. These data have been correlated with the simplified band-theory model that had been proposed earlier. The measurements indicate that the two stimulation bands are associated with the same center. This can be represented by the assignment, to that center, of two energy levels within the energy gap.

SOME OPTICAL PROPERTIES OF n-TYPE InP. R. Newman, General Electric Research Laboratory, Schenectady, New York (Received May 26, 1958).

Measurements of the intrinsic absorption edge of n-type InP at 77°K and 300°K are reported. Differences are found between the spectra of samples of differing origin. The effects are believed due to impurities. Reflection and absorption spectra in the vicinity of the reststrahlen peak are shown. The reststrahlen wavelength is 30.5μ, the static and high-frequency dielectric constants are ϵ_0 =15 and ϵ_∞ =10.6. Free-carrier absorption and reflection spectra are shown. A brief discussion of the implications of the various optical and electrical measurements as regards the conduction band structure is given.

VACANCY PAIRS IN IONIC CRYSTALS. A. B. Lidiard, Department of Physics, University of Reading, Berkshire, England (Received May 20, 1958).

Equations for the thermal equilibrium concentrations of Schottky vacancies and vacancy pairs are presented; these equations allow for the effects produced by Debye-Hückel screening.

AMMONIUM HYDROGEN SULFATE: A NEW FERROELECTRIC WITH LOW COERCIVE FIELD. R. Pepinsky, K. Vedam, S. Hoshino, and Y. Okaya, The Pennsylvania State University, University Park, Pennsylvania (Received May 16, 1958).

$(NH_4)HSO_4$ is ferroelectric in the temperature range between -3°C and -119°C. The symmetry of the room-temperature phase is $P2_1/c$; the ferroelectric phase is also monoclinic, with space group Pc; and the lowest phase, which is piezoelectric, has triclinic symmetry $P1$. All three phases are pseudo-orthorhombic. The coercive field is very low: about 150v/cm at -13°C. The spontaneous polarization P_S is about 0.4 μcoulombs/cm^2 at that temperature, and rises to a maximum of about 0.8 μcoulombs/cm^2 just above -119°C. P_S falls to zero abruptly at the lower transition point. The higher transition is of second order, and the lower of first order.

ROOM-TEMPERATURE FERROELECTRICITY IN LITHIUM HYDRAZINIUM SULFATE, $Li(N_2H_5)SO_4$. R. Pepinsky, K. Vedam, Y. Okaya, and S. Hoshino, Department of Physics, The Pennsylvania State University, University Park, Pennsylvania (Received May 26, 1958).

$Li(N_2H_5)SO_4$ is ferroelectric over the temperature range from below -15°C to above 80°C, and shows no dielectric peaks between 77°K and 140°C. No thermal anomalies are observed. The crystals are orthorhombic, space group $Pna2_1$, with a =8.969A, b=9.913A, c = 5.178A, and the

ferroelectric axis is of course along c; ρ=1.966 g/cc, and there are 4 molecules per cell. The coercive field is about 320 volts/cm at room temperature, and the spontaneous polarization is 0.30 microcoulombs /cm². Beautiful large crystals can be grown from water by a variety of methods. After electroding, the crystal plates should be protected by silicone oil. Without such protection, reaction with normal atmosphere for prolonged time results in excessive conductivity.

NEUTRON-BOMBARDMENT DAMAGE IN SILICON. G. K. Wertheim, Bell Telephone Laboratories, Incorporated, Murray Hill, New Jersey (Received May 22, 1958).

Neutron-bombardment damage in silicon is compared to electron-bombardment effects which have been previously analyzed. A discrete energy level 0.27 ev above the valence band, previously ascribed to the acceptor member of a defect pair having a separation greater than 50A, is found to be produced by both neutrons and electrons. A spectrum of energy levels running from 0.16 ev below the conduction band toward the middle of the gap is ascribed to a defect pair with variable spacing, and is related to the discrete level 0.16 ev below the conduction band which was found previously in electron-bombardment of silicon, and was there ascribed to a closely spaced pair of defects. Lifetime effects are found to be dominated by levels near the middle of the energy gap, which may be related to the above spectrum. A very rapid decrease in mobility at low temperature is ascribed to bombardment-induced inhomogeneities.

PARAMAGNETIC RESONANCE SPECTRUM OF GADOLINIUM IN HYDRATED LANTHANUM TRICHLORIDE. M. Weger and W. Low, Department of Physics, The Hebrew University, Jerusalem (Received May 15, 1958).

The paramagnetic resonance spectrum of Ga^{3+} in $LaCl_3 \cdot 7H_2O$ was measured and found to agree quite well with a spin Hamiltonian with dominant coefficients $b_2^0 = \pm 0.0131\ cm^{-1}$, $b_2^2 = \mp 0.0075\ cm^{-1}$ at room temperature, and $b_2^0 = \pm 0.0099\ cm^{-1}$, $b_2^2 = \mp 0.0115\ cm^{-1}$ at liquid air temperature.

DISSOCIATION OF THE HYDROGEN MOLECULE ION BY ELECTRON IMPACT. Eugene V. Ivash, The University of Texas, Austin, Texas (Received March 11, 1958).

The dissociation cross section of H_2^+ for bombardment by electrons is investigated taking into account the effect of exchange collisions by use of the Born-Oppenheimer approximation. Both the post and prior interactions are considered. Theoretical results are presented for the transition in which the molecule ion is raised from the ground state, $1s\sigma\ ^2\Sigma_g^+$, to the lowest-lying repulsive state, $2p\sigma\ ^1\Sigma_u^+$, and then dissociates. The Franck-Condon principle is assumed throughout, and a classical average is taken over the molecular orientations. Numerical results are given for the dissociation cross sections as a function of energy both with and without the inclusion of exchange effect. The agreement between the post and prior interaction cross sections near the threshold energy is poor.

MAGNETIC MOMENTS DUE TO ROTATION IN Li^6F AND Li^7F. Allan M. Russell,* Syracuse University, Syracuse, New York (Received April 22, 1958).

A molecular beam electric resonance spectrometer has been used to measure the Zeeman splitting of the hyperfine structure of Li^6F and Li^7F in the J=1 rotational state. From these measurements the magnitudes of the rotational magnetic moments of these molecules were determined.

$$(\mu_J/J)_{Li^6F} = 0.0818\ ^{+0.0005}_{-0.0003}\ \text{nuclear magnetons per } J$$

$$(\mu_J/J)_{Li^7F} = 0.0642\ ^{+0.0004}_{-0.0002}\ \text{nuclear magnetons per } J$$

Since the signs of the rotational magnetic moments could not be measured, the sign of the electric dipole moment was not determined experimentally.

A complete expression for the rotational magnetic moment in a diatomic molecule is derived. When the rotational perturbation of the electron motion is neglected, the complete expression is shown to reduce to an equation for the magnetic moment of two singly charged ions. This ion-pair approximation, applied to the molecule

above, indicates that the values for the moments should be positive.

The nonslip electronic contribution, neglected in the ion-pair approximation, is described in terms of a semiclassical model. Numerical evaluation of the nonslip electronic contribution to the rotational magnetic moment of Li^6F and the nuclear J-dependent magnetic field of the fluorine spin-rotation interaction yields results within ten percent of experimental values. The model shows that these two molecular constants, though the results of the same perturbation, arise from different parts of the charge distribution.

*Now at the University of California, Riverside, California.

IONS IN NITROGEN. Mykola Saporoschenko, Department of Physics, Washington University, St. Louis, Missouri (Received May 19, 1958).

N^+, N_2^+, N_3^+, and N_4^+ ions have been identified in nitrogen gas by use of a mass spectrometer operated with the ion source in the pressure range from 10^{-3} mm Hg to 0.6 mm Hg.

Appearance potentials occur at 15.5 ± 0.2 ev for N_2^+, at 15.8 ± 0.3 ev for N_4^+, at 22.1 ± 0.5 ev for N_3^+, and at 24.2 ± 0.4 ev for N^+; N_3^+ ions thus are formed only at much higher electron energies than N_2^+ but still at 2 ev less energy than N^+.

It is believed that the N_4^+ is formed by the process $N_2^+ + N_2 \rightarrow N_4^{+*}$ (excited vibrationally) and the N_3^+ is formed by the process $N_2^{+*} + N_2 \rightarrow N_3^+ + N$ where N_2^{+*} is an excited ion.

N_4^+ and N_3^+ currents decrease with increasing E/p in the source, suggesting that they may be dissociated by molecular impacts. The N_4^+ ions are considerably more readily lost than the N_3^+ ions.

Since the formation of N_3^+ must necessarily release a nitrogen atom, this process constitutes a form of dissociation of N_2 which may account for the value of the dissociation energy of 7.38 ev found by some methods.

VECTOR INTERACTION IN BETA DECAY. Jeremy Bernstein and Robert R. Lewis, Institute for Advanced Study, Princeton, New Jersey (Received May 27, 1958).

Further experimental consequences of a beta-decay interaction generated by a conserved vector current, as proposed by Feynman and Gell-Mann, are studied. Utilizing the strong analogy between the conserved beta-decay current and the conserved charge current, higher multipole corrections to the usual results for allowed transitions are computed, and are related to the corresponding electromagnetic transitions, as was first done by Gell-Mann. The effects calculated include corrections to spectral shapes, beta-gamma directional correlations, and beta-alpha directional correlations. A multipole expansion of the V-A interaction is carried out, including all those terms normally classed as allowed and second forbidden. Applications to several interesting transitions are discussed.

ANALYSIS OF GAMMA-GAMMA ANGULAR CORRELATIONS INVOLVING MULTIPOLE MIXTURES. R. G. Arns and M. L. Wiedenbeck, Harrison M. Randall Laboratory of Phusics, University of Michigan, Ann Arbor, Michigan (Received May 26, 1958).

A method is outlined which provides for rapid analysis of directional correlations involving gamma rays of mixed multipolarity. Specific examples are given to illustrate the application to cascades in which one or both transitions involve a mixture of dipole and quadrupole radiation.

EXTRAPOLATION PROCEDURE FOR THE LOW-ENERGY DEUTERON-DEUTERON REACTION CROSS SECTION. J. G. Brennan, United States Naval Ordnance Laboratory, White Oak, Silver Spring, Maryland, and The Catholic University of America, Washington, D. C. (Received February 6, 1958; revised manuscript received July 8, 1958).

A semiempirical study is made of the deuteron-deuteron reaction cross section to determine its energy dependence at low energy and to permit more reliable extrapolation to energies below the experimental range. An energy-dependent correction to the simple W. K. B. form is shown to be indicated by experimental data. The energy-independent normalization factor is then determined by comparison with the experimental data over the range of ten to one hundred kev bombarding energy. This provides a cross section formula which is more suitable for extrapolation to the region around one kev.

EFFECTS OF VACUUM POLARIZATION SCATTERING IN THE TREATMENT OF PROTON-PROTON SCATTERING DATA. Michiel de Wit and Loyal Durand, III,* Yale University, New Haven, Connecticut (received May 7, 1958; revised manuscript received June 18, 1958).

The effects of including vacuum polarization scattering in angular momentum states with $L>0$ and of employing the relativistic value of the Coulomb scattering parameter η_γ rather than its nonrelativistic value η have been considered in the analysis of low energy p-p scattering data. It has been found that about half of the apparent mean P-wave phase shift present in published data, in the energy range 1.8-4.2 Mev, should be attributed to this vacuum polarization scattering. A further decrease in magnitude of the P wave is obtained when η_γ is employed, rather than η. Both the vacuum polarization scattering in states with $L>0$ and the employment of η_γ rather than η result in a contribution to the f function which does not vary linearly with energy. The effects of these contributions on the coefficients of a polynomial expansion of f in powers of E are ascertained with reference to p-p scattering data in the energy range 0.2 to 7.5 Mev. The object of the work is more that of investigating the possible magnitude of effects on conclusions drawn from data due to applying the various corrections than that of deriving final phase shift values.

*Now at the Institute for Advanced Study, Princeton, New Jersey.

FURTHER STUDIES OF ANOMALOUS INELASTIC PROTON SCATTERING. Bernard L. Cohen and Allen G. Rubin, Oak Ridge National Laboratory, Oak Ridge, Tennessee (Received May 12, 1958).

Anomalous inelastic proton scattering was studied by (1) determining its dependence on bombarding energy, (2) measuring the energy distributions with greatly improved energy resolution, (3) extending the data to separated isotopes and to elements not previously investigated because of target preparation difficulties, and (4) observing the energy distribution of the de-excitation gamma rays following the reaction. In (1), the anomalous peak is found to stay at the same Q value as the bombarding energy is varied. This indicates that the effect is not due to a giant resonance effect, but rather to the regular occurrence of a certain type of level in the final nucleus. In (2), many low-lying levels are resolved; a strong correlation is found between their cross sections for excitation in these reactions and in Coulomb excitation, and an anticorrelation is found between their cross sections for excitation in these reactions and in (p,d) reactions. This is interpreted as good evidence that the reactions used here strongly excite collective levels and only weakly excite single-particle transitions. The fine structure of the anomalous peaks differ greatly among neighboring nuclides, reflecting even-versus-odd and closed shell effects; however, the gross structure as obtained with poor resolution is very similar for neighboring elements. A few striking similarities between the spectra of several elements are pointed out. Several levels of known spin and parity are identified as contributing strongly to the anomalous peak. In (3), it is found that anomalous inelastic scattering ends by becoming irregular and weak between atomic numbers 54 and 64; it is apparently present in every isotope of every element between atomic numbers 30 and 53. In (4), gamma-ray transitions direct to the ground state following anomalous inelastic scattering are found to be relatively common.

Many theoretical explanations for anomalous inelastic scattering are excluded by these results; there is reasonably good evidence that it arises from some type of collective excitation "dissolved" among many shell model states in that energy region.

THE THEORY OF REACTIONS $H^3(d,n)He^4$ AND $He^3(d,p)He^4$. William S. Porter,* Yale University, New Haven, Connecticut; Benjamin Roth,† University of Connecticut, Storrs, Connecticut; and John L. Johnson‡, Yale University, New Haven, Connecticut (Received May 15, 1958).

The reactions $He^3(d,n)He^4$ and $He^3(d,p)He^4$ are fitted using one-level and one-level-plus-background dispersion formulas and taking into account the implications of the principle of charge symmetry of nuclear forces. Fits are found using a variety of penetration parameters and channel radii. The schematic treatment of Breit and the R-matrix formalism of Wigner and Eisenbud are used to obtain the one-level-plus-background formulas and fits made using them are compared with those obtained by previous workers using one-level formulas. The con-

nection of these reaction fits with recent n+He^4 scattering data is discussed.

*Now at the Los Alamos Scientific Laboratory, Los Alamos, New Mexico.
†Now at Oklahoma State University.
‡ Now with the Commercial Atomic Power Authority, Westinghouse Corporation, Pittsburgh, Pennsylvania.

GAMMA-RAY DECAY IN Hg^{200}. R. E. Segel,* Aeronautical Research Laboratory, Wright-Patterson Air Force Base, Dayton, Ohio (Received May 28, 1958).

γ-γ cascades following thermal neutron capture in Hg^{199} have been measured. A three-crystal pair spectrometer was used to measure the high-energy gamma rays and a single NaI(Tl) crystal spectrometer detected the lower energy radiation. A detailed comparison is made between the experiment reported here, the previously known high-resolution mercury capture gamma-ray spectrum, and the previously measured decay scheme of Tl^{200}. Several discrepancies between the capture gamma-ray and β-decay work are resolved, though others remain. By combining the results of the three experiments, probable spin and parity assignments are made for several of the low-lying levels in Hg^{200}. With these spin assignments, some conclusions are drawn as to the relative transition probabilities for the various multipolarity radiations present.

*Guest scientist at Brookhaven National Laboratory, Upton, New York.

MEASUREMENTS OF LIGHT MASSES WITH THE MASS SYNCHROMETER. Lincoln G. Smith,* Brookhaven National Laboratory, Upton, New York (Received May 13, 1958).

A complete account is presented of all results obtained with the mass synchrometer before November, 1957. Included, besides measurements made before March, 1957, tentative values of most of which have been published with little discussion by the author and/or by others, is an extensive later set of measurements on 40 parent ion doublets. From least-squares analyses of 39 of these doublet values and of 11 selected earlier values and from three other earlier values are derived what are believed to be the best values obtainable from all synchrometer data of the mass excesses of the 14 nuclides from which all ions studied are composed. These are H^1, D^2, He^4, B^{10}, B^{11}, C^{12}, N^{14}, O^{18}, Ne^{20}, Ne^{22}, S^{32}, Cl^{35}, Cl^{37}, and A^{40}. Best values of all measured doublets are also tabulated. Statistical analyses indicate that, for external consistency of results, the internal errors given should be multiplied by either 2.35 or 3.6 (the larger factor applying primarily in cases where ions are formed by molecular dissociation). They also indicate that, except in cases where ions are formed with appreciable kinetic energies, the external error of the mean of twenty measurements of a doublet separation made with present techniques is about one part in 4×10^7 of the mass number. The agreement of the best values with the latest comparable values obtained at the University of Minnesota is excellent for the important mass excess C-12 and satisfactory in nearly all other cases. Agreement with comparable values obtained from measurements of nuclear reaction energies is fairly good except for C_3-HCl^{35} and Cl^{35}-35, the synchrometer values of which appear to be in error.

*Now on leave at Project Matterhorn, Princeton, N. J.

RESONANCE THEORY OF NEUTRON CROSS SECTIONS OF FISSIONABLE NUCLEI. Erich Vogt, Atomic Energy of Canada Limited, Chalk River, Ontario, Canada (Received May 15, 1958)

The general Wigner-Eisenbud resonance theory is used to develop a method of analysis for the neutron cross sections of fissionable nuclei. The method is employed in giving a reasonable description of the low energy cross sections in U^{235}. The single-level fit for U^{235} is known to be unreasonable. Many-level expressions for the cross sections are derived, the only approximation to the general theory being the neglect of all but a small group of resonances. No explicit reference to fission channels is needed and the many-level expressions require few level parameters: the E_λ, $\Gamma_{\lambda n}$, $\Gamma_{\lambda\gamma}$, and $\Gamma_{\lambda f}$ of the single-level theory for each resonance and a few additional parameters pertinent to the interference between levels. The interference terms are described and shown to be important. Their average value yields information about the number of channels involved in fission. The shape and size of the U^{235} cross sections below 2 ev are fitted to within one or two per cent using (a) only one negative-energy resonance of smaller size than in

the single-level fit; (b) no additional levels to fit the shape other than the observed levels at positive energies; (c) three interference parameters whose size suggests that there are several fission channels in U^{235}.

SLOW-NEUTRON RESONANCES IN U^{235}. F. J. Shore and V. L. Sailor, Brookhaven National Laboratory, Upton, New York (Received May 15, 1958).

The total and fission cross sections of U^{235} have been measured in the energy range 0.1 to 10 ev on the BNL crystal spectrometer. The capture cross section is derived by combining these data with published scattering cross section data. Analysis of the low-energy resonance structure depends strongly on the interpretation of the "background" cross section which dominates the low-energy region. This "background" cannot be accounted for in terms of observed resonances; however, if it is attributed to the presence of a bound level, approximate parameters for such a level can be obtained. It is seen that the strength Γ_n^0 of the hypothetical bound level is anomalously large.

The capture component of the resonances is found to be symmetrical, i.e., of Breit-Wigner shape. However, the fission component of the same resonances is asymmetrical, which indicates that a multilevel formula is necessary for fitting the resonances in the fission cross section. A good but not perfect fit to the data between 0 and 2 ev has been achieved using a multilevel formula which applies to the case of a single fission channel and many capture channels. The remaining small discrepancies in fitting the data indicate that the assumption of a single fission channel is too restrictive, and that one or more additional channels contribute weakly to slow neutron fission in U^{235}.

BETA DECAY OF N^{16}. David E. Alburger, Brookhaven National Laboratory, Upton, New York (Received May 29, 1958).

The positron-electron pair coincidence spectrum occurring in the beta decay of 7.4-sec N^{16} has been examined at 1.25% resolution with an intermediate-image pair spectrometer. A line is observed corresponding to internal pair conversion of the 6.14-Mev $E3$ transition from the 3- second excited state to the 0+ ground state of O^{16}. The upper limit on the intensity of a 6.06-Mev nuclear pair line from the 0+ first excited state (electric monopole transition) is 10% relative to the 6.14-Mev line and this corresponds to an upper limit of 1.5×10^{-4} for the fractional beta decay to the 6.06-Mev level. A lower limit of 8.2 is derived for the log ft value of this 2- → 0+ beta-ray branch as compared with the known log ft = 6.7 for the 2- → 0+ 10.4-Mev branch to the ground state.

RADIATIVE CAPTURE OF PROTONS IN OXYGEN AT 140 TO 170 kev. R. E. Hester, R. E. Pixley, and W. A. S. Lamb, University of California Radiation Laboratory, Livermore, California (Received May 26, 1958).

The thick-target yield of the reaction $O^{16}(p,\gamma)F^{17}$ has been measured by bombarding Al_2O_3 targets with protons from 140 to 170 kev using currents from 3 ma to 10 ma and counting the induced positron activity of the F^{17}. The thick-target yield ranges from $(1.9\pm1)\times10^{-17}$ beta/incident proton at 140 kev to $(1.47\pm0.15)\times10^{-16}$ beta/incident proton at 170 kev. The corresponding cross sections are $(4.3\pm2.4)\times10^{-11}$ barn at 140 kev and $(2.34\pm0.3)\times10^{-10}$ barn at 170 kev. The activity was identified by observing the half-life. The cross section factor S_0 was found to be 6.8 ± 1.4 kev-barns between 140 kev and 170 kev bombarding energy.

THERMAL NEUTRON INDUCED FISSION OF Th^{229}. A.Smith, P. Fields, A. Friedman, and R. Sjoblom, Argonne National Laboratory, Lemont, Illinois (Received May 26, 1958).

A back-to-back ionization chamber is used to measure the Th^{229} fission fragment kinetic energy distribution. This fission process is found to be the most asymmetric known, with a most probable mass ratio of 1.57. The measured total fragment energy is 160.2 ± 3 Mev. The experimental results are compared with the characteristics of other fission processes and with theory.

π^--p INTERACTIONS AT 1.85 Bev/c. R. C. Whitten and M. M. Block, Duke University, Durham, North Carolina (Received May 29, 1958).

This paper reports some diffusion cloud chamber results concerning the elastic and inelastic

scattering in hydrogen of pions having a momentum in the laboratory system of 1.85 Bev/c. The elastic scattering data are consistent with diffraction by a sphere of radius 0.85×10^{-13} cm and opacity 0.9. The forward scattering amplitude was found to be in agreement with that derived from the dispersion relations. Investigation of the inelastic scattering disclosed that there is little agreement with the predictions of the Fermi statistical model. The angular and Q-value distributions of the particles emerging from inelastic interactions showed little agreement with either the statistical theory or the excited isobar model.

PHOTON DECAY OF HYPERONS. R. E. Behrends, Brookhaven National Laboratory, Upton, New York (Received April 14, 1958; revised manuscript received May 23, 1958).

The photon mode of decay of the Σ^+ and Λ^0 is examined. The lifetime and angular distributions of the decay products are written in terms of the three physical parameters of the problem without restriction as to invariance under space reflection, charge conjugation, or time reversal. An estimate of these parameters is made by use of perturbation theory, and, although ambiguous, it indicates that this mode may soon be experimentally detectable.

EFFECT OF RECOIL ON THE ELASTIC SCATTERING OF HIGH-ENERGY ELECTRONS BY ZERO-SPIN NUCLEI. Leslie L. Foldy, Case Institute of Technology, Cleveland, Ohio, and Kenneth W. Ford, Benny J. Hill, David L. Hill, and John G. Wills, Los Alamos Scientific Laboratory, University of California, Los Alamos, New Mexico (Received August 12, 1958).

The effect of nuclear recoil on the elastic scattering of high-energy electrons or muons by zero-spin nuclei is taken into account in exact phase-shift calculations by adapting the Breit two-particle Hamiltonian to the case that one of the two particles is of finite size, is spinless, and is nonrelativistic, the other being a normal point Dirac particle. A radial and angular separation of the Dirac equation is still possible. Illustrative exact calculations demonstrate the recoil effect on the scattering of 420-Mev electrons by carbon, oxygen, and lead. At large angles, the dynamic recoil effect is substantially greater than the kinetic recoil effect. However, conclusions about the shape and size of the nuclear charge distribution are not significantly altered.

PION PRODUCTION IN PION-NUCLEON COLLISIONS AT MODERATE ENERGIES. Saul Barshay, Physics Department, Brookhaven National Laboratory, Upton, New York (Received May 7, 1958).

Some remarks are made concerning pion production in pion-nucleon collisions at incident kinetic energies of from 300 Mev to about 600 Mev. We discuss a certain representation of the matrix element which involves explicitly a term that is dependent upon a pion-pion scattering amplitude off the energy shell. A partial wave decomposition of the rest of the transition operator is attempted upon the assumption that the amplitudes so introduced contain the effect of strong pion-nucleon interactions in the P wave. We discuss certain quantities, such as the angular distribution and the polarization of the recoil nucleon, which may be determined largely by the interference between the amplitudes involving specific pion-pion and pion-nucleon effects, respectively.

EMPIRICAL MODEL FOR ULTRARELATIVISTIC NUCLEON-NUCLEON COLLISIONS. Giuseppe Cocconi, Laboratory of Nuclear Studies, Cornell University, Ithaca, New York (Received May 29, 1958).

The interactions of ultrarelativistic ($\geq10^{12}$ ev) primary cosmic-ray protons with nucleons in photographic plates observed thus far show that in the majority of the cases the secondaries are emitted isotropically from two centers; this fact and the observation that the average transverse momentum of the secondaries is independent both of the energy of the primary nucleon and of the energy of the secondaries themselves suggest the following empirical model. When a nucleon-nucleon collision takes place at these energies, two "bodies" are formed, that move in opposite directions with respect to the center of momentum of the system, together with the two original nucleons, stripped of the energy necessary to create the "bodies". Each "body" emits, in its own system of reference, about half of the total number of secondaries, each secondary having an average energy of ~1 Bev, independent of its nature. The inelasticity of the

collision is thus determined by the number of secondaries and by the velocity of each "body" in the center of momentum of the two original nucleons. Accurate determinations of the energy, in the laboratory system, of the secondaries of these interactions could provide a sensitive test for the model, but at present there are too few measurements available.

CHARGE INDEPENDENCE AND THE LOW-ENERGY PARAMETERS IN PION PHYSICS. Michael J. Moravcsik, University of California Radiation Laboratory, Livermore, California (Received May 26, 1958).

It is pointed out that the striking discrepancy among the very-low-energy parameters in pion physics might be resolved by considering deviations from charge independence. An accurate set of measurements of the total cross section for charge exchange scattering at low energies is needed to prove or disprove this contention.

Σ^0 MASS. M. Lynn Stevenson, Radiation Laboratory, University of California, Berkeley, California (Received May 22, 1958).

The mass of the Σ^0 hyperon is obtained from a single associated production event in which both strange particles are observed to decay. The event is unusual, first because the Σ^0 hyperon decays directly into an electron-positron pair plus a Λ, and secondly because all the tracks are long and of higher-than-average quality. The value of Q_Σ obtained from this event is 76.4 ± 3.3 Mev where $Q_\Sigma \equiv M_{\Sigma^0} - M_\Lambda$. For $M_\Lambda = 1115.2 \pm 0.14$ Mev the mass of the Σ^0 is 1191.6 ± 3.3 Mev. Upon averaging this value with the existing measurements of Alvarez et al. and Eisler et al., one obtains $Q_\Sigma = 75.3_{-1.4}^{+0.9}$ Mev and hence $M_{\Sigma^0} = 1190.5_{-1.4}^{+0.9}$ Mev.

FORM FACTORS IN QUANTUM ELECTRODYNAMICS. S. D. Drell and F. Zachariasen, Stanford University, Stanford, California (Received May 19, 1958).

The electromagnetic form factors of an electron in pure quantum electrodynamics (QED) are analyzed with the techniques of dispersion relations. The viewpoint is adopted here that no subtractions are required in the construction of dispersion relations for the electromagnetic vertex. This leads to coupled integral equations for the form factors in terms of other physical amplitudes; electron-positron scattering, for example. The relation between this and the usual perturbation approach to QED, and the validity and consequences of the "no subtraction" philosophy, are discussed.

FOURTH-ORDER NUCLEON-NUCLEON DISPERSION RELATIONS. M. T. Grisaru, Palmer Physical Laboratory, Princeton University, Princeton, New Jersey (Received May 21, 1958).

The fourth-order contribution to the nucleon-nucleon scattering amplitude is shown to satisfy the conditions necessary for writing dispersion relations. The imaginary part is explicitly calculated and shown to agree with the absorptive part obtained from general principles.

PHOTOMESON PRODUCTION IN INTERMEDIATE COUPLING. Kurt Haller,* Department of Physics, Columbia University, New York, New York (Received November 21, 1957).

The wave functions for the ground state of a physical nucleon and for the scattering state of a physical nucleon and a pion are used to compute photopion production cross sections in an intermediate coupling approximation. Since the mesons included in the physical nucleon suffice to account for an approximately correct value of the anomalous magnetic moment, no use is made in this theory of a correction term involving experimental values of the latter. The p-wave part of the photoproduction matrix element is arrived at naturally, largely from the interaction of the meson currents with the electromagnetic field. Differential cross sections are obtained for the production of π^+ and π^0 by photons of 200, 260, 335, and 440 Mev in the laboratory system. It is also pointed out that in the intermediate coupling theory the ambiguity in defining a gauge-invariant Hamiltonian is nontrivial. Different procedures for generalizing the meson-nucleon interaction term into a form that is gauge invariant in an extended-source theory lead to substantially different p-wave components of the photoproduction cross sections.

* Present Address: Wayman Crow Laboratory of Physics, Washington University, St. Louis, Missouri.